Doctoral Thesis

Hole transport and spin-orbit coupling in p-type GaAs nanostructures

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Hole transport and spin-orbit coupling in p-type GaAs nanostructures

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presented by

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Abstract

In this thesis we investigate hole transport and spin-orbit coupling effects in p-type GaAs nanostructures. The interest in p-type GaAs structures arises primarily from the fact that spin-orbit interactions are much stronger in this material system compared to the well explored n-type GaAs structures. Therefore, p-type GaAs nanostructures are potentially important for the development of spintronics. We have developed the fabrication of p-type GaAs nanodevices and observed evidence for ballistic, phase-coherent and single-hole transport in these structures.

We explore a novel type of p-type GaAs/AlGaAs heterostructures obtained by doping with carbon. The high quality of the C-doped GaAs heterostructures is demonstrated by the observation of the Fractional and Integer Quantum Hall effect, as well as by highly resolved Shubnikov-de Haas (SdH) oscillations in low temperature magnetotransport experiments. The mobility anisotropy in these 2DHGs grown in (100) direction is found to be much smaller compared to the mobility anisotropy in previously studied, Si-doped (311) p-type GaAs heterostructures. The mean free path in these 2DHGs is larger than 1.5 µm at a temperature of 70 mK, which allows the fabrication of ballistic p-type nanostructures.

The simultaneous observation of weak antilocalization and the beating in SdH oscillations documents the presence of strong spin-orbit (SO) interactions in C-doped p-type GaAs heterostructures. A Fourier analysis of the beating of the SdH oscillations allows us to determine the populations of the two spin-split subbands, which originate from the SO induced splitting of the heavy hole band. From the charge imbalance between these two subbands we estimate the SO induced splitting $\Delta_{SO}$ of the heavy hole subband to be around 30% of the total Fermi energy for both investigated heterostructures, which proves the presence of exceptionally strong SO interactions. Weak antilocalization measurements allow us to quantify phase coherence in these p-type GaAs heterostructures. The measured phase coherence length $l_{\phi}$ of holes at a temperature of 70 mK is larger than 1.8 µm, which is promising for the fabrication of phase-coherent p-type nanodevices.

In order to make nanostructures on p-type GaAs we employ a local anodic oxidation technique with an Atomic Force Microscope (AFM). This approach enables us to overcome the problems with hysteresis and sample instabilities, present in devices fabricated with conventional split-gate techniques. Using AFM oxidation lithography we have fabricated a quantum point contact (QPC), a quantum dot.
and two quantum rings. All nanostructures are fabricated from the same p-type GaAs heterostructure with a 2DHG located 45 nm below the sample surface.

The QPC conductance reveals clear plateaus at $2e^2/h$ and $4e^2/h$ which confirms the ballistic nature of hole transport through the QPC. Besides, the observation of the additional plateau-like structure around $0.8 \times 2e^2/h$ points to the signature of the so-called 0.7 feature in our QPC. We observe the zero-bias peak in the source-drain bias spectroscopy for the QPC set around $0.8 \times 2e^2/h$, which might indicate a Kondo-like origin of the 0.7 feature.

We observe pronounced and reproducible Coulomb blockade resonances in transport measurements performed on a single hole transistor. This is the first observation of Coulomb blockade in a p-type GaAs quantum dot. A charging energy of around 1.5 meV is extracted from Coulomb diamond measurements, in agreement with the lithographic dimensions of the dot. Due to the large effective mass of the holes, the estimated single-particle level spacing for this quantum dot (around 15 µeV) is of the order of the thermal broadening, and therefore it was not possible to resolve excited states in the differential conductance measurements. This absence of excited states as well as the temperature dependence of Coulomb peak heights indicate that the dot is in the multi-level transport regime. Nevertheless, fluctuations in peak spacings larger than the estimated mean single-particle level spacing are observed which could be related to stronger carrier-carrier Coulomb interactions in p-type systems. The observation of Coulomb blockade effects in another sample with different geometry confirms the ability to reproducibly fabricate electronically functional p-type GaAs quantum dots with AFM lithography.

Phase coherent transport in two quantum ring samples is further explored. We have measured Aharonov-Bohm (AB) oscillations with a visibility of around 3% in a larger ring with an orbital radius of 420 nm fabricated on a p-type GaAs heterostructure with exceptionally strong spin-orbit interactions. Clear beating patterns observed in the raw data can be interpreted in terms of a spin geometric phase. Beside the h/e oscillations, we resolve the contributions from the second harmonic of AB oscillations and also find a beating in these h/2e oscillations. A resistance minimum at $B = 0$ T, present in all gate configurations, is the signature of destructive interference of the spins propagating along time-reversed paths. The AB oscillations persist up to 350 mK. From their temperature dependence we estimate the phase-coherence length of holes to be around 2 µm at the base temperature of 60 mK, compatible with weak antilocalization measurements. In the smaller ring with the orbital radius of 160 nm we have measured AB oscillations with a visibility larger than 10%. Minimum at $B = 0$ T in all gate configurations is again observed, in agreement with results from the large ring sample.
Zusammenfassung


Wir untersuchen einen neuartigen Typ von p-GaAs/AlGaAs Heterostrukturen, der durch die Dotierung mit Kohlenstoff (C) zustande kommt. Die hohe Qualität von C-dotierten GaAs Heterostrukturen zeigt sich in der Beobachtung des fraktionierten und ganzzahligen Quanten-Hall-Effekts ebenso wie in den hoch aufgelösten Shubnikov-de Haas (SdH) Oszillationen in Magnetotransport-Experimenten bei tiefer Temperatur. Die Mobilitäts-Anisotropie in diesen in (100)-Richtung gewachsenen zwei-dimensionalen Loch-Gasen (2DHGs) ist viel kleiner verglichen mit der Mobilitäts-Anisotropie in Si-dotierten (311) p-GaAs Heterostrukturen früherer Studien. Die mittlere freie Weglänge in diesen 2DHGs ist größer als 1.5 µm bei T = 70 mK.


Um Nanostrukturen auf p-GaAs herzustellen, verwenden wir eine lokale Oxidations Technik mit einem Rasterkraftmikroskop (AFM). Diese Vorgehensweise ermöglicht uns, Probleme mit Hysterese und Proben-Instabilitäten bei lateralen Gate-Elektroden zu bewältigen. Mit AFM-Oxidations-Lithographie stellten wir
einen Quanten-Punkt-Kontakt (QPC), einen Quanten-Dot und zwei Quanten-Ringe her. Alle Proben stammen von der gleichen p-GaAs Heterostruktur.

Die Leitfähigkeit des QPC zeigt deutliche Plateaus bei $2e^2/h$ und $4e^2/h$, was die ballistische Natur des Löchertransports durch die QPC bestätigt. Ausserdem deutet die Beobachtung von zusätzlichen, plateau-ähnlichen Strukturen um $0.8 \times 2e^2/h$ in unserem QPC auf die Signatur der so genannten 0.7 Struktur hin. In der Leitwertspekroskopie der Zu- und Ableitung des QPC beobachten wir um $0.8 \times 2e^2/h$ ein Maximum bei Vorspannung Null, was möglicherweise den Kondo-artigen Ursprung dieser 0.7 Struktur andeutet.


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<th>explanation</th>
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<tr>
<td>$a_B$</td>
<td>Bohr radius</td>
</tr>
<tr>
<td>$-e&lt;0$</td>
<td>electron charge</td>
</tr>
<tr>
<td>$\epsilon$</td>
<td>dielectric permittivity</td>
</tr>
<tr>
<td>$\epsilon_0$</td>
<td>vacuum dielectric constant</td>
</tr>
<tr>
<td>$h = 2\pi \hbar$</td>
<td>Planck’s constant</td>
</tr>
<tr>
<td>$k_B$</td>
<td>Boltzmann constant</td>
</tr>
<tr>
<td>$\mu_B$</td>
<td>Bohr magneton</td>
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</tbody>
</table>

<table>
<thead>
<tr>
<th>Abbreviation</th>
<th>Explanation</th>
</tr>
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<tr>
<td>2DHG</td>
<td>two dimensional hole gas</td>
</tr>
<tr>
<td>2DEG</td>
<td>two dimensional electron gas</td>
</tr>
<tr>
<td>AB</td>
<td>Aharonov-Bohm</td>
</tr>
<tr>
<td>AAS</td>
<td>Altshuler-Aronov-Spivak</td>
</tr>
<tr>
<td>ac</td>
<td>alternating current</td>
</tr>
<tr>
<td>AFM</td>
<td>atomic force microscope</td>
</tr>
<tr>
<td>CB</td>
<td>Coulomb blockade</td>
</tr>
<tr>
<td>dc</td>
<td>direct current</td>
</tr>
<tr>
<td>DOS</td>
<td>density of states</td>
</tr>
<tr>
<td>FWHM</td>
<td>full width at half maximum</td>
</tr>
<tr>
<td>LL</td>
<td>Landau level</td>
</tr>
<tr>
<td>MBE</td>
<td>molecular beam epitaxy</td>
</tr>
<tr>
<td>PG</td>
<td>plunger gate</td>
</tr>
<tr>
<td>QD</td>
<td>quantum dot</td>
</tr>
<tr>
<td>QHE</td>
<td>quantum Hall effect</td>
</tr>
<tr>
<td>QPC</td>
<td>quantum point contact</td>
</tr>
<tr>
<td>SdH</td>
<td>Shubnikov de Haas</td>
</tr>
<tr>
<td>SO</td>
<td>spin-orbit</td>
</tr>
<tr>
<td>TG</td>
<td>top gate</td>
</tr>
<tr>
<td>Symbol</td>
<td>Explanation</td>
</tr>
<tr>
<td>--------</td>
<td>-------------</td>
</tr>
<tr>
<td>L,W</td>
<td>system size (length, width)</td>
</tr>
<tr>
<td>$B$</td>
<td>magnetic field</td>
</tr>
<tr>
<td>$\vec{A}$</td>
<td>vector potential</td>
</tr>
<tr>
<td>$C_\Sigma$</td>
<td>self-capacitance of a quantum ring/dot</td>
</tr>
<tr>
<td>$E_C$</td>
<td>constant interaction energy</td>
</tr>
<tr>
<td>$\alpha_G$</td>
<td>gate lever arm</td>
</tr>
<tr>
<td>$\varepsilon_N$</td>
<td>single particle energy of the Nth level</td>
</tr>
<tr>
<td>$E_F$</td>
<td>Fermi energy</td>
</tr>
<tr>
<td>$\Delta_{SO}$</td>
<td>spin-orbit induced splitting of the valence band</td>
</tr>
<tr>
<td>$Q$</td>
<td>charge</td>
</tr>
<tr>
<td>$I$</td>
<td>current</td>
</tr>
<tr>
<td>$V$</td>
<td>voltage</td>
</tr>
<tr>
<td>$G$</td>
<td>conductance</td>
</tr>
<tr>
<td>$R$</td>
<td>Resistance</td>
</tr>
<tr>
<td>$\rho$</td>
<td>resistivity</td>
</tr>
<tr>
<td>$D$</td>
<td>density of states</td>
</tr>
<tr>
<td>$\Gamma^{S,D}$</td>
<td>tunnel coupling</td>
</tr>
<tr>
<td>$v_F$</td>
<td>Fermi velocity</td>
</tr>
<tr>
<td>$k_F$</td>
<td>Fermi wavevector</td>
</tr>
<tr>
<td>$\lambda_F$</td>
<td>Fermi wavelength</td>
</tr>
<tr>
<td>$\ell_c$</td>
<td>magnetic length</td>
</tr>
<tr>
<td>$\ell_{e, m}$</td>
<td>elastic mean free path</td>
</tr>
<tr>
<td>$\ell_\phi$</td>
<td>phase coherence length</td>
</tr>
<tr>
<td>$\tau_{tr}$</td>
<td>transport scattering time</td>
</tr>
<tr>
<td>$\tau_\phi$</td>
<td>phase coherence time</td>
</tr>
<tr>
<td>$\tau_q$</td>
<td>quantum life time</td>
</tr>
<tr>
<td>$\tau_{SO}$</td>
<td>spin-orbit scattering time</td>
</tr>
<tr>
<td>$m^*$</td>
<td>effective mass</td>
</tr>
<tr>
<td>$\mu$</td>
<td>mobility</td>
</tr>
<tr>
<td>$N$</td>
<td>charge sheet density</td>
</tr>
<tr>
<td>$\nu$</td>
<td>filling factor</td>
</tr>
<tr>
<td>$\omega_c$</td>
<td>cyclotron frequency</td>
</tr>
<tr>
<td>$r_c$</td>
<td>cyclotron radius</td>
</tr>
<tr>
<td>$r_s$</td>
<td>interaction parameter</td>
</tr>
<tr>
<td>$g^*$</td>
<td>effective g-factor</td>
</tr>
<tr>
<td>$s_z$</td>
<td>z - component of the spin</td>
</tr>
<tr>
<td>$T$</td>
<td>temperature</td>
</tr>
</tbody>
</table>

ix
Chapter 1

Introduction

"A hole?" the rock chewer grunted. "No, not a hole", said the will-o’-the-wisp despairingly. "A hole, after all, is something. This is nothing at all".

Michael Ende - "The Neverending Story"

Two theories which revolutionized physics and science in general at the beginning of the twentieth century are Quantum Mechanics and Theory of Relativity. While the first describes the behavior of ultra-small object, the second deals with very fast objects whose velocity approaches the speed of light. One of the phenomena which originates from the interplay between quantum and relativistic effects is the coupling between spin and orbital degrees of freedom of electrons in atoms.

On the other side progress in the physics of semiconductors in the second half of the twentieth century played a key role for the development of modern technologies. Very few discoveries in the history of science had such an enormous impact on the development of mankind as the discovery of the transistor by Bardeen, Brattain and Shockley in 1947 [1]. There are billions of transistors in almost every electronic device around us - according to Moore’s law [2] the number of transistors in integrated circuit doubles every 18 months. This miniaturization of electronic components will reach the level where their functionality is affected by quantum effects.

However, modern electronics is still based only on manipulating the charge of electrons. It is well established that besides its electric charge the electron has a magnetic moment - the spin. Therefore, the natural question which arises is if it is possible to utilize the electron spin instead of its charge as a carrier of information. This emerging research field whose objective is to manipulate spins in solid state devices is called spintronics. Advantages of spintronics compared to present charge-based electronics should be nonvolatility, less electric power consumption and higher data processing speeds [3]. Conventionally, spin manipulation is realized with external magnetic fields, produced by large magnetic coils. Such an approach faces the problem of nonscalability and is not suitable for on-chip realization. On the
other side, the possibility to manipulate spins electrically would allow using present state-of-the-art semiconductor fabrication technologies.

Spin-orbit coupling represents the major link which would allow spin manipulation with electric fields. Therefore, the exploration of spin-orbit interaction effects in low-dimensional semiconducting systems received considerable attention recently, since the detailed understanding and control of spin-orbit coupling in solid-state devices is crucial for the development of spintronics. The basic concept of a spintronics device, proposed by Datta and Das [4], is a spin field-effect transistor. The modulation of the current between source and drain in the proposed device arises from spin precession due to spin-orbit coupling controlled by an external gate.

Besides being significant for progress in spintronics, systems with strong spin-orbit coupling provide also important experimental ground for studying some of the fundamental quantum mechanical concepts, such as geometrical phases [5, 6]. The Aharonov-Bohm phase [7], which represents the geometric phase acquired by the orbital wave function of the charged particle encircling a magnetic flux line is experimentally well established and manifests itself through the oscillations in the resistance of mesoscopic rings in an external magnetic field. The Aharonov-Casher phase [8] represents the geometric phase acquired by the spin upon its evolution along a closed path in the spin’s Hilbert space. The spin’s geometric phase has been observed in experiments with neutrons [9], but its observation in solid state systems appeared to be challenging and is presently the subject of intensive experimental work [10–14]. It is theoretically predicted that the spin’s geometric phase modulates the transport properties of mesoscopic ring with a spin-orbit induced orientationally inhomogeneous magnetic field [15–20]. We observe clear beating patterns in Aharonov-Bohm oscillations measured in p-type GaAs mesoscopic rings which results in a splitting of $h/e$ and $h/2e$ peaks in the Fourier spectrum. These results present additional indications of the spin’s geometric phase in solid state systems.

The strength of spin-orbit interaction, however, varies greatly in different semiconductors. In the two most extensively studied semiconductors, Si and n-type GaAs, spin-orbit interactions are relatively weak and therefore the investigation of spin-orbit interactions is directed towards exploring new materials such as p-type GaAs, InAs, HgTe, InSb or AlSb. Thus, besides exploring new physical phenomena emerging in these materials, a significant part of research in this field is still devoted to the development of nanofabrication technologies and optimization of material properties.

In this thesis we investigate low-temperature hole transport and spin orbit interaction effects in novel p-type GaAs structures, obtained by doping with carbon. Although the concept of holes in solid state physics is one of the major achievements of the semiclassical theory of electron dynamics in solids, the attitude of physicists towards semiconductor valence bands remained somewhat reluctant, mainly due to the more complex structure of the valence bands. However, instead of considering these peculiarities of the GaAs valence band as a difficulty, one may look at them as a source of new, potentially interesting phenomena.
The dispersion relation of the GaAs valence band is such that holes have several times larger effective masses than electrons - thus the Fermi energies in two dimensional hole systems are much lower compared to the conventional two dimensional electron systems with comparable densities and interaction corrections become more expressed in hole systems. Consequently, low-dimensional hole systems can serve as a very fertile ground for exploring not only spin-orbit coupling effects, but also carrier-carrier Coulomb interactions. With typical values of the interaction parameter (ratio between Coulomb and Fermi energy) $r_s > 5$, our two-dimensional hole gases enter the strongly interacting regime.

The second important characteristic of the valence band is that the states at the top of the valence band have p-like symmetry (orbital angular momentum $l = 1$), while those at the bottom of the conduction band have s-like symmetry ($l = 0$). This larger orbital angular momentum of holes is another reason for the stronger spin-orbit interaction in hole systems compared to electron doped GaAs systems. The difference in the symmetry of the hole (p) and electron (s) wave functions has also an important impact on the strength of the nuclear hyperfine interaction in hole and electron systems. Due to the p-like symmetry of the hole states, the overlap of their wave functions with nuclear wave functions is much smaller compared to the overlap of electron and nuclear wave functions. Consequently, the nuclear hyperfine interaction, which represents one of the main sources of spin decoherence is expected to be much weaker in holes compared to electron systems.

Research on electronic transport in low-dimensional p-type GaAs systems was, so far, mainly limited to two-dimensional bulk samples. Phenomena like the metal-insulator transition, the fractional quantum Hall effect, and the formation of different electronic phases in strongly interacting regimes are intensively studied in two-dimensional hole systems over the last years. However, the investigation of hole transport in one and zero-dimensional p-type GaAs nanodevices received significantly less attention, due to technological difficulties to fabricate stable p-type GaAs nanostructures with conventional split-gate techniques [21]. In order to overcome these problems with metallic gates, we employ in this thesis a different lithography technique, namely, Atomic Force Microscope (AFM) oxidation lithography [22, 23] to define tunable nanostructures on two-dimensional hole gases. The electronic functionality of these structures was demonstrated by observing conductance quantization in a quantum point contact, and for the first time Coulomb blockade in p-type GaAs quantum dot.

The thesis starts with a short introduction into the basic physics concepts required to understand the experimental results presented in this work (chapter 2). In chapter 3 we discuss the technology of sample fabrication with emphasis on the advantages of the AFM defined samples compared to conventional split-gate defined p-type structures. Chapter 4 is devoted to the technology of making Au/Zn/In Ohmic contacts and presents the effects related to the fact that InZn alloys may become superconducting in the milikelvin temperature range where the experiments are performed. The high quality of the carbon doped p-type GaAs heterostructures
is demonstrated by observing highly resolved Shubnikov-de Haas oscillations as well as fractional quantum Hall states by low-temperature magnetotransport experiments (chapter 5). In the same chapter the evidence for strong spin-orbit interaction is provided by the simultaneous observation of a beating in Shubnikov-de Haas oscillations and classical positive magnetoresistance. Chapter 6 presents weak anti-localization measurements in the limit of very strong spin-orbit interactions. The phase coherence time of holes is extracted from the fit of the weak anti-localization peak. Conductance quantization and the so called '0.7 feature' in a p-type quantum point contact are discussed in chapter 7. In chapter 8 we present Coulomb blockade measurements in two p-type GaAs quantum dots with different geometries. Finally, in chapter 9 Aharonov-Bohm oscillations in the presence of strong spin-orbit interactions are explored in two quantum rings with different sizes and the obtained results strongly point to the signatures of the spin's geometric phase.
Chapter 2
Basic concepts

In this chapter we introduce general concepts necessary for the understanding of the experiments presented in this thesis. Detailed discussions of more specific concepts related to the presented experiments, as well as comparisons of our experiments to those previously reported in literature will be given separately at the beginning of the chapter related to the particular experiment.

2.1 Electronic transport in low-dimensional semiconductor systems

2.1.1 Two dimensional electron (hole) gas

A two-dimensional electron (hole) gas is a system containing charge carriers which are strongly confined in one direction and free to move in the two other directions. Charge carriers are confined in a quantum well formed at the interface of a semiconductor heterostructure (for technological details see paragraph 3.1). When all charge carriers reside in the lowest quantized state of a quantum well the system of charge carriers behaves effectively as two-dimensional. The typical spatial extension of a 2D system in the confinement direction is around 10 nm. Realization of 2D electron (hole) gases allowed discoveries of several fascinating physics phenomena like the Integer and Fractional Quantum Hall effects [24, 25] (Nobel Prizes in 1985 and 1998). Besides, 2D electron (hole) gases represent the ground for the fabrication of low-dimensional nanostructures and the exploration of quantum effects present in such systems.

Two-dimensional electron (hole) gases have a peculiar property that their spin degenerate density of states is independent of energy:

\[
\mathcal{D} = \frac{m^*}{\pi \hbar^2}
\]  

(2.1)

where \(m^*\) is the effective mass of a charge carrier. Expressions for the Fermi energy and the Fermi wavelength follow directly from the above formula:
Chapter 2. Basic concepts

\[ E_F = \frac{\pi \hbar^2}{m^* N}, \quad \lambda_F = \sqrt{\frac{2\pi}{N}} \]  

(2.2)

where \( N \) is the sheet density of the charge carriers. The Fermi wavelength \( \lambda_F \) is around 40 nm in the heterostructures used in this thesis. When the size of the nanostructures is of the order of the Fermi wavelength quantum confinement effects become important.

### 2.1.2 Magnetotransport in two-dimensional systems

Two-dimensional electron (hole) systems are typically characterized by performing low-temperature magnetotransport experiments on a Hall-bar geometry (Fig. 2.1(a)). The two main characteristics of such systems are their density \( N \) and mobility \( \mu \), which can be extracted using the following relations from the classical Drude theory of magnetotransport:

\[ N = \frac{1}{e(d\rho_{xx}/dB)|_{B=0}} \quad \mu = \frac{1}{eN\rho_{xx}(B = 0T)}, \]  

(2.3)

where \( \rho_{xx} \) and \( \rho_{xy} \) are the measured longitudinal and transversal resistivity, respectively.

The classical Drude theory of conductivity is valid only for weak magnetic fields, for which \( \omega_c \tau \ll 1 \), with \( \omega_c = eB/m^* \) being the cyclotron frequency, and \( \tau \) the elastic scattering time. At high enough B-fields where the previous relation does not hold anymore, the constant density of states (eq. (2.1)), present at zero magnetic field, evolves into a series of \( \delta \)-peaks, the so-called Landau levels (LLs), evenly spaced in energy by \( \hbar \omega_c \). In real samples these \( \delta \)-peaks are broadened into Gaussian-like peaks due to scattering off impurities. Since the separation between the Landau levels and the degeneracy of each level increase at higher B-fields, the number of LLs below the Fermi level, the so-called filling factor \( \nu \), decreases. Because the total number of carriers in the system is constant, the number of carriers per LL (the so-called degeneracy) increases with magnetic field and reads: \( N_L = eB/h \).

The two effects which originate from such a formation of Landau Levels in 2D systems are the Shubnikov-de Haas (SdH) oscillations and the integer Quantum Hall effect \([26, 27]\). When the filling factor \( \nu \) is integer, the Fermi level lies exactly between the two LLs and the density of states has a minimum, leading to a minimum in both, longitudinal resistivity \( \rho_{xx} \) and conductivity \( \sigma_{xx} \) \([26, 27]\). At low enough temperatures and high enough magnetic fields, when the LLs are well developed, the resistance at the positions of the SdH minima is exactly zero. At the same B-fields where minima in \( \rho_{xx} \) develop, the Hall resistivity \( \rho_{xy} \) displays quantized plateaus at precisely:

\[ \rho_{xy} = \frac{\hbar}{e^2 \nu}, \nu = 1, 2, 3, ... \]  

(2.4)
The examples of minima in Shubnikov de Haas oscillations and plateaus in the integer quantum Hall effect are shown in Fig. 2.1(b) for filling factors $\nu = 1, 2, 3, 4$.

The explanation of the integer quantum Hall effect is based on the concept of edge states. When the Fermi level lies between two LLs, where the density of states within the sample is zero, the current is carried only along the sample edges via the so-called edge-states. Landau levels which are below the Fermi level in the bulk sample intersect the Fermi level at the sample edges and form edge states. Therefore, the number of edge states is equal to the filling factor $\nu$. The transport along these edge states is dissipationless and ballistic, and therefore the Landauer-Büttiker formalism [29, 30] can be employed for the calculation of the Hall resistivity ((eq. (2.4))) at the plateaus. The values of the resistance at the quantized plateaus are determined with extreme precision due to the fact that edge states which propagate in opposite directions are along opposite edges of the Hall
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bar. This macroscopic separation of counterpropagating edge states completely suppresses backscattering, leading simultaneously to zero values of $\rho_{xx}$ and precisely quantized values of $\rho_{xy}$.

We can see in Fig. 2.1(b) that beside at the integer values of the filling factor $\nu$, minima in $\rho_{xx}$ and plateaus in $\rho_{xy}$ can also occur at filling factors of the form $\nu = p/q$, where $p$ is integer and $q$ is odd integer number. This effect is called Fractional Quantum Hall effect [25] and is present in 2D samples with exceptionally high mobility. The Hall resistance at the quantized plateaus is in that case given as

$$\rho_{xy} = \frac{h}{e^2} \frac{1}{p/q},$$

(2.5)

The fractional quantum Hall effect originates from strong carrier-carrier Coulomb interactions and its theoretical description can be given in the framework of many-body wavefunctions. However, the introduction of new quasiparticles, called composite fermions, which consist of one electron and two attached flux quanta, allows the description of the fractional quantum Hall effect in the framework of non-interacting composite fermions. In this picture the fractional quantum Hall effect can be seen as the integer quantum Hall effect of composite fermions [27].

The observation of the fractional quantum Hall effect demonstrates at the same time the presence of strong Coulomb interactions and the high quality of the sample.

2.1.3 Conductance quantization in quantum point contacts

It is well known that the conductance of a metallic wire, whose dimensions are much larger then the electron mean free path, decreases continuously as its cross-section shrinks. But, what happens when the lateral dimensions of the wire become smaller than the average distance between impurities in the material? Advances in nanofabrication technologies make the answer to this question possible.

A quantum point contact is a narrow constriction defined in a 2D electron (hole) system whose width $W$ is of the order of the Fermi wavelength of the charge carriers and much smaller than the mean free path in the system. The conductance of such a system was first measured by van Wees et al. [31] (Fig 2.2) and Wharam et al. [32] and it was found that, rather than decreasing continuously with a narrowing constriction, the QPC conductance decreases in quantized steps of $2e^2/h$.

Landauer and Büttiker introduced a new concept of conductance in ballistic systems [29, 30]. They directly relate the conductance of a ballistic device to its transmission:

$$G = \frac{2e^2}{h} T = \frac{2e^2}{h} \sum_{n,m=1}^{N} |t_{nm}|^2,$$

(2.6)

where $T$ is the total transmission, $t$ is the transmission matrix and summing is performed over all transverse modes in the QPC. In the case of ideally transmitted
2.1. Electronic transport in low-dimensional semiconductor systems

one-dimensional modes the conductance of a QPC is given as

$$G = \frac{2e^2}{h} N,$$  \hspace{1cm} (2.7)

where $N$ is an integer number. This expression can be understood as each transverse mode within a QPC carries a conductance quantum of $2e^2/h$. The number of transverse modes within a QPC is given as $N = 2W/\lambda_F$, where $\lambda_F$ is the Fermi wavelength. As the electronic width $W$ of a QPC decreases, the number of transverse modes transmitted through the QPC also decreases, leading to the step-wise decrease of the QPC conductance (Fig. 2.2). Beside plateaus at integer multiples of $2e^2/h$ very often an additional plateau-like structure around $0.7 \times 2e^2/h$ is observed. Experimental observations related to this so-called 0.7 structure and its possible origins will be discussed in more detail in paragraph 6.1.

It should be mentioned that the Landauer-Büttiker formalism can be generalized to multi-terminal ballistic devices.

2.1.4 Coulomb blockade in quantum dots

Quantum dots are small conducting islands, that can have a tunable number of confined charge carriers. Due to this spatial confinement of charge carriers in all three dimensions the energy spectrum of the confined carriers is discrete. The two characteristic energy-scales in quantum dots are the charging energy, which represents the energy necessary to add one more charge carrier to the dot, and the quantum mechanical confinement energy.

In this thesis we perform spectroscopy of energy spectra of p-type GaAs quantum dots by performing low-temperature electronic transport measurements. The dot is
coupled to the macroscopic current and voltage leads (source and drain) through tunnel barriers, which can be well controlled with external gate voltages. If the size of the dot, and therefore its electrostatic capacitance is small enough that the charging energy $E_C$, necessary to add one more charge carrier to the dot is larger than the thermal excitation energy $k_BT$, transport of charge carriers through the dot is blocked and the so-called Coulomb blockade effect occurs.

A detail review on electronic transport in quantum dots can be found in [33]. Here we mention briefly the basic concepts necessary for understanding of our measurements on p-type quantum dots presented in chapter 8.

A scheme of a quantum dot tunnel-coupled to source and drain contacts and capacitively coupled to the gate is shown in Fig. 2.3(a). The total electrostatic energy of the dot presented in Fig. 2.3(a) is given by

$$U(N) = \frac{(Ne - C_g V_g)^2}{2C}$$

where $C$ is the total capacitance of the dot, $C_g$ is the gate capacitance and $Ne - C_g V_g$ is the total charge of the dot, which consists of the $N$ electrons on the
dot and displacement charge $C_g V_g$ which is induced by the gate, and can be varied continuously.

The energy necessary to add the $N^{th}$ electron to the dot is therefore

$$
\mu_{\text{dot}}(N) = U(N) - U(N - 1) = \frac{e^2}{C} (N - \frac{1}{2}) - e C_g V_g
$$

(2.9)

When this energy lies within the source-drain bias window $\mu_S \geq \mu_{\text{dot}}(N) \geq \mu_D$ (see Fig. 2.3(c), lower-left), electron transport through the dot can occur through the so-called sequential tunnelling process. Otherwise, when $\mu_{\text{dot}}$ is outside of the source-drain bias window (Fig. 2.3(c), lower-right) transport through the dot is blocked. It should be mentioned that in the experiments, the source-drain bias window is typically chosen to be very small $\mu_S \simeq \mu_D$, so that only single energy levels in the dot can be probed in transport.

The addition energy, which corresponds to the change of the electrochemical potential needed to increase the number of electrons on a dot by one at a fixed gate voltage $V_g$ is given as:

$$
\Delta \mu = \mu(N + 1) - \mu(N) = \frac{e^2}{C}.
$$

(2.10)

However, at fixed electrochemical potentials of the source and drain $\mu_S \simeq \mu_D$ this addition energy has to be compensated by the change of a gate voltage $\Delta V_g$ in order for transport through the dot to occur. Thus, $\mu(N, V_g) = \mu(N + 1, V_g + \Delta V_g)$. This condition can be satisfied for any number of electrons on a dot for a suitably chosen gate voltage and therefore the whole series of Coulomb peaks develops in the dot conductance. The gate-voltage separation between the neighboring peaks is given by

$$
\Delta V_g = \frac{1}{e \alpha} \frac{e^2}{C}
$$

(2.11)

where $\alpha = C_g/C$ is the gate lever arm.

Beside the charging energy, the quantum-mechanical confinement energy in semiconductor quantum dots might also play an important role. In the constant interaction model \[35\], the quantum mechanical level spacing $\Delta E = E_N - E_{N-1}$ simply adds to the charging energy in order to give the total addition energy of the $N^{th}$ electron:

$$
\Delta \mu(N) = \frac{e^2}{C} + \Delta E_N
$$

(2.12)

If $\Delta E \gg k_B T$ the dot is in the so-called single-particle level regime, and transport occurs via single quantum level. On the other side for $\Delta E \ll k_B T$, the dot is in the multi-level transport regime and the transport occurs via many single-particle levels. In the case of a dot with steep potential walls, the mean single-particle level spacing can be calculated as $\Delta = 2\pi \hbar^2/g m^* A$, where $g$ is the degeneracy of electron
(hole) states and \( A \) is the electronic area of the dot. Due to the fact that holes in 2D systems have a much larger effective mass than electrons, the mean single-particle level spacing in the case of hole quantum dots is much smaller, and therefore harder to resolve, compared to electron quantum dots.

A sufficiently large source-drain bias voltage can also lift the Coulomb blockade. The electrostatic considerations for the finite biases applied across the dot show that stability regions in \((V_g, V_{sd})\) parameter space, where transport through the dot is blocked, are diamond-shaped regions (Fig. 2.3(c)). Within these diamond-shaped regions the number of electrons in the dot is a fixed integer number. At the edges of the diamonds the electron number can fluctuate and conduction sets in. The diamonds intersect the \( V_g \) axis at gate voltages where the Coulomb peaks are observed.

### 2.1.5 Aharonov-Bohm and Altshuler-Aronov-Spivak oscillations

The wave nature of corpuscular objects was the ground-breaking concept which quantum mechanics introduced into modern physics. Interference phenomena with particles represent the most distinct manifestation of the particles’ wave nature. In order to observe the particles’ interference effects in transport properties of mesoscopic conductors it is important that the characteristic length over which particles (electrons or holes) preserve the information about their phase (the so-called phase-coherence length \( l_\phi \)) is larger than the characteristic geometric size of the investigated system.

The Aharonov-Bohm effect \[7\] is an interference effect based on a phase a charge particle acquires upon traversing a ring like mesoscopic structure threaded by a magnetic flux line. The scheme of the interfering trajectories which cause the Aharonov-Bohm effect is shown in Fig. 2.4. A phase difference acquired by the charge particle propagating along the two arms of the ring \( \gamma_1 \) and \( \gamma_2 \) in the presence of the external magnetic vector potential \( \vec{A} \) is given by

\[
\delta(\phi) = \delta(0) - \frac{e}{\hbar} \int_{\gamma_1 - \gamma_2} \vec{A} d\vec{s} = \delta(0) - 2\pi \frac{\phi}{\phi_0},
\]

(2.13)

where \( \delta(0) \) is the phase difference between the two paths in the absence of external flux, which depends on the microscopic configuration of the ring, \( \phi \) is the external magnetic flux, and \( \phi_0 = \hbar/e \) is a flux quantum. The transmission through such a ring-like interferometer is given by \[27\]

\[
T = T_1 + T_2 + 2\sqrt{T_1 T_2} \cos(\delta(0) - 2\pi \frac{\phi}{\phi_0}),
\]

(2.14)

where \( T_1 \) and \( T_2 \) are the individual transmissions through the upper or lower arm, respectively. We can see that the total transmission through the ring oscillates as a function of the external magnetic flux. According to the Landauer-Büttiker
formalism [29, 30], the conductance of a ballistic mesoscopic device is proportional to its transmission, and therefore the ring conductance also oscillates as a function of the external magnetic flux. These $h/e$ periodic oscillations were first observed in a solid state device by Webb et al. [36] in a small metallic ring, and their first observation in a ring fabricated from semiconductor heterostructure was made by Timp et al. [37].

![Figure 2.4: (a) Scheme of the interfering trajectories that cause $h/e$ Aharonov-Bohm oscillations](image)

However, the two interference paths shown in Fig. 2.4 are not the only paths which contribute to interference. The paths with larger number of windings around the ring also interfere and give rise to higher harmonics of the AB oscillations. But due to the larger length of these additional paths, the visibility of higher harmonics becomes smaller. Particularly important is the second harmonic of the AB oscillations which originates from the interference of the time reversed paths which encircle the ring exactly once - these $h/2e$ oscillations are called Altshuler-Aronov-Spivak (AAS) oscillations [38]. The importance of these $h/2e$ oscillations comes from the fact that the phase difference for zero external flux, $δ(0)$, between the two time-reversed paths is zero, and due to that AAS $h/2e$ oscillations are much more robust than $h/e$ oscillations when the microscopic configuration of the ring is changed. Therefore, $h/2e$ oscillations persist even when averaging over a whole ensemble of microscopic ring configurations is performed. This is the reason why $h/2e$ oscillations were observed before $h/e$ AB oscillations in thin metallic cylinders [39].

It was later shown that the Aharonov-Bohm phase is a special case of the more general geometric phase [5, 6] acquired by the orbital wave function of a charged particle encircling a magnetic flux line. This so-called Berry’s phase [5] has a purely geometric origin and depends exclusively on the topology of the parameter space of the quantum system, and the path in this parameter space along which the quantum system performs an adiabatic cyclic evolution [40]. This concept was generalized also for the non-adiabatic cyclic evolution of a quantum system in a parameter space [6].

Aharonov-Bohm oscillations, as mentioned above, originate from the phase ac-
quired by the orbital part of electron’s (hole’s) wavefunction upon its cyclic propagation in a ring. However, charge carriers possess also a spin-degree of freedom, which can acquire an additional spin-related geometric phase in the presence of spin-orbit interactions. This spin-related geometric phase and its consequences on Aharonov-Bohm oscillations are discussed in detail in paragraph 9.1.

2.1.6 Weak localization and antilocalization

Weak localization is a quantum mechanical effect which arises from the constructive interference between time reversed partial waves of charge carriers in disordered materials. This constructive interference leads to an enhanced probability of backscattering of carriers and therefore the longitudinal resistance of the samples increases. The application of a perpendicular magnetic field breaks time reversal symmetry and suppresses weak localization, which manifests as an appearance of a resistance peak in the longitudinal magnetoresistance at B=0T.

![Figure 2.5: Scheme of the time-reversed trajectories which interfere constructively and produce weak localization.](image)

The weak localization effect has exactly the same origin as the Altshuler-Aronov-Spivak $\hbar/2e$ oscillations, with the only difference that instead of a well-defined ring loop, the whole ensemble of time-reversed loops, formed by the elastic scatterers (Fig. 2.5) contribute to backscattering. The contributions from all these pairs of loops have to be summed up, with proper weights, related to their areas. The smallest loops which contribute to the sum have the size of the order of $l_m^2$ ($l_m$ is the mean free path) and the largest have the size of the order of $l_\varphi^2$ ($l_\varphi$ is the phase coherent length). Each pair of the time reversed loops gives AAS oscillations with the phase $\varphi = 0$ at $B = 0$T, and the frequency related to the area of the loop. Larger loops give higher oscillation frequencies. The AAS oscillations will therefore have all possible phases at $B \neq 0$T and only at $B = 0$T, they will all be in phase ($\varphi = 0$). Thus, the sum over the whole ensemble of loops produces a resistance peak at $B = 0$T and averages out the oscillations at all other B-fields. The width of the weak localization resistance peak is determined by the value of the B-field at which the slowest oscillations, related to the loop with area $l_m^2$, have a phase $\pi$.
2.2 Spin-orbit coupling effects in GaAs valence band

2.2.1 Origin of spin-orbit interaction

Spin-orbit interaction was first investigated in atomic physics in the context of the fine structure of atoms. It represents the relativistic correction of the atomic energy spectra. Spin-orbit correction to the “hydrogen-like” Hamiltonian of atoms emerges from the expansion of the relativistic Dirac equation into the non-relativistic Pauli equation and reads:

$$H_{SO} = \frac{\hbar}{4e^2m^2} (\nabla V(\vec{r})) \times \vec{p} \cdot \vec{\sigma},$$  \hspace{1cm} (2.15)

where $V$ is the Coulomb potential of the atomic core, $\vec{p}$ is the momentum operator and $\vec{\sigma} = (\sigma_x, \sigma_y, \sigma_z)$ is the vector of the Pauli spin matrices. The expression eq. (2.15) can be transformed into $H_{SO} = \xi(r) \vec{l} \cdot \vec{s}$, where $\vec{l}$ is the electron’s orbital angular momentum, $\vec{s}$ is spin angular momentum and $\xi(r)$ is a scalar function of a radial variable $r$. This relation clearly shows how the orbital and spin angular momentum are coupled in atoms. Spin-orbit coupling scales with the atomic number of the atom and therefore is more important for heavier elements. Due to this dependence, spin-orbit interactions are weak in Si, compared to GaAs, InAs or InSb.

The relativistic origin of the spin-orbit interaction can be visualized in the following way: in the reference frame where the atom is at rest, only the electric field from the core of the atom is present. However, the electron which orbits around the atomic core, sees in its own rest frame beside the electric, also the magnetic field given by

$$\vec{B} = \frac{1}{c^2} \vec{\sigma} \times \vec{E},$$  \hspace{1cm} (2.16)

due to relativistic transformations of the fields. This magnetic field in the rest frame of the electron couples to the electron spin via the Zeeman interaction.
### 2.2.2 Valence band of bulk GaAs

The motion of electrons in crystals is characterized by energy bands $E_n(\vec{k})$, with band index $n$ and wavevector $\vec{k}$. A schematic representation of the GaAs band structure close to the $\Gamma$ point is shown in Fig. 2.6(a), while the calculated GaAs bands using the pseudopotential method [42] are shown in Fig. 2.6(b). GaAs is a direct semiconductor with an energy gap around 1.42 eV. It has zinc-blende crystal structure.

![Figure 2.6](image-url)

**Figure 2.6:** (a) Schematics of the GaAs bands structure close to the $\Gamma$ point with conduction band, heavy and light hole valence band and spin-orbit split-off band shown. (b) Band structure of GaAs calculated using the pseudopotential method including spin-orbit effects [42].

The conduction band has $s$-like symmetry and its Bloch states are characterized by an orbital quantum number $l = 0$. On the other hand states at the top of the valence band have $p$-like symmetry and their orbital quantum number is $l = 1$. This is one of the reasons why spin-orbit effects are stronger for holes than for electrons. The total angular momentum of holes $J = l + s$ ($s$ is the spin quantum number) can therefore have two values $J = 3/2$ and $J = 1/2$. Due to spin-orbit interaction this sextet of bands is split into a doublet with $J = 1/2$ (split-off SO bands) and the quadruplet with $J = 3/2$ (heavy (HH) and light (LH) holes). The splitting between the split-off bands and a quadruplet of $J = 3/2$ hole bands at the $\Gamma$ point is 0.34 eV. The four hole bands with $J = 3/2$ are degenerate at the $\Gamma$ point ($\vec{k} = 0$) but for non-zero wavevectors this degeneracy is lifted and two doublets with $J_z = \pm 3/2$ and $J_z = \pm 1/2$ are formed. The doublet with $J_z = \pm 3/2$ is called heavy-hole (HH) band and the doublet with $J_z = \pm 1/2$ is called light-hole (LH) band. The names heavy and light hole band come from the fact that the effective masses of these two sets of bands are different. The heavy hole band has an effective mass $m_{hh} = 0.53m_0$. 

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2.2. Spin-orbit coupling effects in GaAs valence band

and the light hole $m_{lh} = 0.08m_0$, where $m_0$ is free electron mass. However, it should be mentioned here that in contrast to the conduction band which is parabolic and can be well described with an effective mass approximation, heavy and light hole bands are highly non-parabolic and anisotropic, and the simple description with an effective mass approximation is not satisfactory. Therefore, the hole effective masses depend on the orientation and magnitude of the wavevector, and the values quoted above give only rough estimates. It should still be emphasized that the effective mass of heavy holes $m_{hh} = 0.53m_0$ is significantly larger than the effective mass of electrons $m_e = 0.067m_0$.

2.2.3 Two-dimensional hole gases

The nanostructures investigated in this thesis are fabricated from two-dimensional hole gases in GaAs heterostructures. We discuss in chapter 3 how such a confinement of holes into two dimensions is technologically achieved. In this paragraph we focus on the analysis of hole states in 2DHGs. The size quantization of holes in the heterostructure’s growth directions has a significant impact on the relation between the hole states in a 2DHG to that in bulk GaAs. The fourfold degeneracy of $J = 3/2$ hole bands bands (HH,LH) in the center of the Brillouin zone is lifted by size quantization, and the HH states with $J_z = \pm 3/2$, and the LH states with $J_z = \pm 1/2$ become separated in energy at $k_{\parallel} = 0$ as shown in Fig. 2.7(a) ($k_{\parallel} = 0$ is the in-plane wavevector in the case of a 2DHG). Subband quantization induces HH-LH mixing, and therefore the anticrossing between HH and LH occurs at finite in-plane wavevector (Fig. 2.7(a)). It should be mentioned that for 2D hole systems the HH (LH) subband has a larger (smaller) effective mass in the growth direction. However, from Fig. 2.7(a) we can see that in the plane of the 2DHG the so-called mass-inversion occurs and therefore the HH subband has smaller in-plane effective mass than the LH subband.

Fig. 2.7(b) shows numerically calculated hole subbands in a 2DHG formed in a 15nm wide quantum well grown in (001) direction [43]. The 2DHG dispersion relations are presented for (100) and (110) in-plane directions. Size quantization of holes in the growth direction produces the whole series of 2D subbands, originating from bulk heavy and light hole subbands. The separation of the bulk split-off subband from the bulk heavy and light hole subbands is large enough so that the split-off subband does not play any role in 2D hole systems. It should be noted in Fig. 2.7(b) that the separation between the topmost thso subbands (HH1 and LH1) is of the order of 10 meV. For typical doping densities $\sim 10^{11}$ cm$^{-2}$ in p-type GaAs heterostructures, the Fermi energy is rather small, due to the large effective mass of holes, typically of the order 2-3 meV. Therefore it is clear that in the quantum limit of low temperatures only the top-most heavy-hole subband is occupied in 2D hole gases.

It is important to emphasize that size quantization of holes in the growth ($z$) direction implies that the quantization axis in the system is along that direction.
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Figure 2.7: (taken from [43]) (a) Lifting of the degeneracy between 2D heavy and light holes and HH-LH anticrosing. (b) Numerical calculation [43] of 2D hole dispersion relations for a GaAs heterostructure grown along (001) direction. Dispersions are calculated for (100) and (110) in-plane directions.

This fact has important consequences on heavy holes in quantum wells since they have an orbital quantum number $l = 1$, which is quantized along the $z$-direction. Since there can exist only one quantization axis in the system, the hole’s spin has also to be aligned along the $z$-direction and therefore heavy holes behave as particles with effective spin $J = 3/2$ which can have projections on the $z$-axis $J_z = \pm 3/2$. Therefore, the Zeeman interaction for in-plane B-fields or spin-orbit interaction, which tend to align the hole spins in the plane of the 2D hole system, have to compete with the energy of the HH-LH splitting which enforces the quantization axis along the growth directions. This is the reason why, in the case of heavy holes in the 2DHG, the Zeeman splitting for an in-plane magnetic field is a higher order effect and the in-plane g-factor of holes which gives a linear contribution to Zeeman energy is $g = 0$ [43]. Similarly, this is also the reason why spin-orbit interactions in 2D hole systems which tend to orient the spins in the plane are characterized by cubic in wavevector splitting, and not linear, like in the case of electrons.

In the end of this paragraph it should be mentioned that due to the significantly larger effective mass of holes compared to electrons in 2D systems (see paragraph 5.6), Fermi energies in 2D hole systems are much lower than in electron systems with comparable densities, which make 2D hole gases particularly suitable for the investigation of interaction effects (spin-orbit or carrier-carrier Coulomb interactions).
2.2. Spin-orbit coupling effects in GaAs valence band

2.2.4 Spin-orbit coupling in two-dimensional electron and hole systems

The spin degeneracy of energy bands in solid state systems, $E_{k\uparrow} = E_{k\downarrow}$, results from the spatial inversion symmetry of the crystal lattice and time inversion symmetry. The lack of the spatial inversion symmetry in a solid leads to a lifting of the spin degeneracy even in the absence of the external magnetic fields. Energy bands in a spatially asymmetric solid-state system possess the so-called Kramers degeneracy $E_{k\uparrow} = E_{-k\downarrow}$.

The lack of inversion asymmetry may have twofold origin: the bulk inversion asymmetry (Dresselhaus SO splitting [44]) or the structural inversion asymmetry (Rashba SO splitting [45]).

The Dresselhaus SO splitting is present in materials whose crystal lattice lacks the center of inversion. The examples of such materials are the III-V semiconductor compounds with zinc-blende structure, such as GaAs or InSb. In bulk GaAs the Dresselhaus SO splitting of the conduction band is cubic in wavevector $k$. The Dresselhaus SO splitting is also present in low-dimensional semiconductor systems which are hosted in materials with bulk inversion asymmetry. The Dresselhaus splitting in 2D electron gases in GaAs is linear in wavevector $k$ and can be expressed as

$$H_D = \beta_D (\sigma_x k_x - \sigma_y k_y),$$

(2.17)

where $\sigma_x, \sigma_y$ are the Pauli spin matrices, and $\beta_D$ is the Dresselhaus coefficient. It depends on band structure parameters of the host material and also on the thickness of the electron gas in the growth direction, $\beta_D = \beta \langle k_z^2 \rangle$. It follows that narrower quantum wells have larger Dresselhaus coefficients. Spin orientation in the plane of a 2DEG in the presence of Dresselhaus SO splitting can be seen in Fig. 2.8(c).

With the advancement of MBE growth technologies it became possible to grow layered semiconductor heterostructures with tailored structural inversion asymmetry. This structural inversion asymmetry induces spin splitting (Rashba splitting) in the 2D electron or hole systems which are hosted in these heterostructures. Rashba SO splitting in 2D electron gases in GaAs is also linear in wavevector $k$ and is given by the following expression

$$H_R = \alpha_R (\sigma_x k_y - \sigma_y k_x),$$

(2.18)

where $\alpha_R$ is the Rashba spin-orbit coupling parameter. This coefficient depends on the electric field $E_z$ perpendicular to the 2D electron gas and can be expressed as $\alpha_R = \alpha \langle E_z \rangle$, where $\alpha$ is a material specific constant and $\langle E_z \rangle$ is the average electric field in z-direction. Electric field $E_z$, and therefore the Rashba SO splitting can be tuned by applying voltages to external gates, which makes the Rashba SO mechanism particularly important for applications in spintronics. Experimentally, the tuning of the SO splitting by external gates was first demonstrated by Engels et
Figure 2.8: (taken from [46]) Scheme of the dispersions of electrons in a 2DEG in the presence of SO interaction. (a) Dispersions for only Rashba, or only Dresselhaus SO splitting. (b) Fermi contours and spin orientations for Rashba term. (c) Fermi contours and spin orientations for Dresselhaus term. (d) Dispersion relations if both Rashba and Dresselhaus terms are present. (e) Fermi contours and spin orientations if both Rashba and Dresselhaus terms are present.

al. [47] and Nitta et al. [48]. Dispersion relations and spin orientations in 2DEGs when Rashba or Dresselhaus or both these SO mechanisms are present, are shown in Fig. 2.8.

SO splitting due to inversion asymmetry in 2D hole gases differs from that in 2D electron gases due to the fact that heavy holes are effectively spin 3/2 particles. It was mentioned in the end of the previous paragraph that due to the size quantization of holes in the growth direction, the quantization axis in the system is enforced to be in the growth direction. Due to that SO interactions, which tend to orient spins in the plane of the 2DHG, represent higher order corrections. Therefore, both Rashba and Dresselhaus SO splittings in 2D hole gases with an occupied heavy-hole band are cubic in wavevector $k$ [43, 49]. The exact Hamiltonian for Rashba SO interaction for heavy holes can be found in [50], and the hole spin orientations resulting from this Hamiltonian displays more complex pattern than that shown for electrons in Fig. 2.8.

The presence of strong SO interactions in 2D hole gases induces a spin-splitting
of the heavy hole band (see Fig. 5.1). This spin-splitting and its experimental manifestations are discussed in detail in paragraph 5.1. The importance of SO interactions can be evaluated by comparing SO induced splitting of the two subbands $\Delta_{SO}$ at the Fermi level with the Fermi energy in the system $E_F$. 
Chapter 3

Sample fabrication and measurement setup

3.1 Carbon doped p-type GaAs/AlGaAs heterostructures

The development of Molecular Beam Epitaxy (MBE) allows growth of layered semiconductor structures with atomic precision. Such layered structures are called heterostructures. This possibility to grow different materials atomic layer by atomic layer provides a way to tailor the band-structure of heterostructures in the growth direction. GaAs/AlGaAs heterostructures show superior quality compared to the other material systems due to the almost perfect lattice match between GaAs and AlGaAs, which ensures almost no growth dislocations from strain relaxation. The scheme of the p-type GaAs/AlGaAs heterostructure that has been studied, together with its bandstructure, calculated by self-consistent Poisson-Schrödinger solver [51], is shown in Fig 3.1. Due to the different band gaps of GaAs and AlGaAs the potential well is formed at the GaAs/AlGaAs interface (inset in Fig.3.1(b)). Doping in the heterostructure shown below is adjusted such that there is only one state due to confinement in z-direction above the Fermi level. Therefore, holes in this state are confined in the z-direction (the width of the confined region is around 10 nm) and are free to move in x and y directions, behaving effectively as a two-dimensional (2D) system.

Another concept, important for the high quality of the GaAs/AlGaAs heterostructures is modulation doping. Employing this technique allows to spatially separate the dopants from the carriers in an underlying 2D system (Fig. 3.1). Due to this spatial separation dopant scattering is suppressed and the mobility of carriers in the 2D system can reach extremely high values, which is important for the observation of various quantum phenomena. In GaAs 2D electron systems the mobility of the electrons can exceed $10^7$ cm$^2$/Vs [52]. Silicon serves as a standard donor in these high mobility n-type GaAs heterostructures, which are typically grown in (100).
direction. Due to such exceptional electronic properties of n-type GaAs heterostructures, most of the transport experiments have been done with GaAs electrons so far.

Figure 3.1: Left: Layout of the heterostructure Bochum12029 with a 2DHG 45 nm below the surface. Right: Profile of the valence band energy along the growth direction (red line) with designated position of the Fermi level (dashed line) calculated by a self-consistent Poisson-Schrödinger solver [51]. Inset: Magnified confining potential of the 2DHG (red line), with the quantized level of the 2DHG (dark blue line) and Fermi level (dashed line). The light blue line shows the spatial profile of the charge density.

The presence of much stronger spin-orbit interactions in p-type compared to n-type GaAs heterostructures, makes GaAs holes promising for the development of spintronics devices. However, the growth of p-type GaAs heterostructures is not as mature as the growth of n-type GaAs heterostructures. Modulation doping in p-type GaAs/AlGaAs heterostructures was introduced with beryllium (Be) as a dopant on (100) structures [53]. Subsequently the integer [54] and the fractional [55] quantum Hall effects were observed in p-type material and transistor action in such devices was demonstrated [56]. Although the mobility could be improved to several $10^4$ cm$^2$/Vs [57], the further development of such devices remained limited due to the high diffusion constant of Be dopants in GaAs. Progress in the growth of high mobility p-type GaAs/AlGaAs heterostructures was made by using silicon (Si) as an acceptor. Si, which has diffusion constant 100 times smaller than Be [58], can be incorporated on arsenic sites in (311)A plane of GaAs, and therefore acts as an acceptor. Two-dimensional hole gases (2DHGs) obtained in such way can reach mobilities of up to $1.2 \times 10^6$ cm$^2$/Vs [59] and the integer and the fractional quantum Hall effects were observed in such systems as well [60]. However, Si doped (311)A 2DHGs also have several drawbacks. The reduced symmetry of the (311)A crystallographic plane results in a more complicated band structure than in the case of (100) 2DHGs. Besides, a surface corrugation with a periodicity of 3.2 nm and a height of up to 5 monolayers, characteristic for (311)A GaAs planes [61] is responsible for a strong mobility anisotropy in (311)A 2DHGs. Typically, the
mobility along the (-233) in-plane direction is 2-4 times larger than the mobility along the (0-11) in-plane direction [62], which complicates the interpretation of transport measurements in these systems. Finally, measurements by Simmons et al. [63] suggest that interface roughness scattering, arising from the nature of the (311)A GaAs surface, is the dominant scattering mechanism which limits a further increase of mobility above $10^6 \text{cm}^2/{\text{Vs}}$ in Si doped (311)A GaAs heterostructures.

Recently it has become possible to dope GaAs with carbon (C) acting as an acceptor on (100) substrates [64]. Carbon, similarly as Si, has a diffusion constant in GaAs about two orders of magnitude smaller than that of Be [58]. The advantage of C over Si is that it acts as an acceptor on the technologically more important GaAs (100) surface, producing more isotropic 2DHGs with a simpler band structure [43].

We have measured for the first time the integer and the fractional quantum Hall effect in C-doped p-type GaAs/AlGaAs heterostructures with a mobility of 160'000 cm$^2/{\text{Vs}}$ [65], demonstrating the high electronic quality of C-doped GaAs/AlGaAs heterostructures. Later, further progress in the fabrication of high mobility C-doped GaAs/AlGaAs heterostructures was made by Gerl et al. [66] and Manfra et al. [67] who simultaneously reported 2DHG mobilities above $10^6 \text{cm}^2/{\text{Vs}}$. Currently, the highest mobility reported for C-doped 2DHG is $2.2 \times 10^6 \text{cm}^2/{\text{Vs}}$ [68], which already exceeds the highest value reported for Si-doped p-type GaAs heterostructures.

We present in the following table the characteristics of the C-doped wafers investigated in this thesis: in the first column is the wafer’s name, in the second the depth of 2DHG below the sample surface, in the third is the growth sequence and in the fourth the measured density and mobility at 4.2K. The notation ‘(Al)GaAs:C doped’ means that the corresponding layer is spatially uniformly doped with a density of C-dopants in the range of $10^{18} \text{cm}^{-3}$. The shallow wafer Bochum12029 with a 2DHG 45 nm below the surface has excellent electronic properties and is suitable for AFM oxidation lithography. The layout of this wafer, together with its valence band profile is shown in Fig. 3.1. All nanostructures measured in this thesis (chapters 7,8,9,10) are fabricated on this wafer. Besides, the Hall-bar fabricated from this wafer is used for Shubnikov-de Haas measurements (chapter 5) and weak antilocalization measurements (chapter 6) in a dilution fridge. The AFM oxidation lithography on another, cap-doped shallow wafer, Bochum11673, was problematic, with oxide lines having a lot of topographic defects, presumably due to doping of the cap layer. The three wafers with deep hole gases (100 nm, 120 nm and 135 nm) are not suitable for fabrication of nanostructures with AFM oxidation lithography. The Hall-bar fabricated from the wafer Bochum1282, with a 2DHG 100 nm below the sample surface, is used for Shubnikov-de Haas and quantum Hall effect measurements (chapter 5) and weak antilocalization measurements (chapter 6) in a dilution fridge. The other two wafers with a deep 2DHG, Bochum11678 and Bochum11680, although showing signatures of Shubnikov-de Haas oscillations at 1.7K, are only used for tests of metallic top-gates presented later in this chapter.

We have further explored the mobility anisotropy in three of these wafers (Bochum12029,
3.1. Carbon doped p-type GaAs/AlGaAs heterostructures

<table>
<thead>
<tr>
<th>wafer</th>
<th>2DHG depth</th>
<th>growth sequence</th>
<th>(n) and (\mu) at 4.2K</th>
</tr>
</thead>
<tbody>
<tr>
<td>Bochum12029</td>
<td>45nm</td>
<td>5nm GaAs</td>
<td>(n = 3.8 \times 10^{11} \text{ cm}^{-2})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>15nm Al_{0.35}Ga_{0.65}As:C doped</td>
<td>(\mu = 120'000 \text{ cm}^2/\text{Vs})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>25nm Al_{0.35}Ga_{0.65}As</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>65nm GaAs</td>
<td></td>
</tr>
<tr>
<td>Bochum1282</td>
<td>100nm</td>
<td>5nm GaAs:C doped</td>
<td>(n = 3.1 \times 10^{11} \text{ cm}^{-2})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>65nm Al_{0.31}Ga_{0.69}As:C doped</td>
<td>(\mu = 130'000 \text{ cm}^2/\text{Vs})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>30nm Al_{0.31}Ga_{0.69}As</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>650nm GaAs</td>
<td></td>
</tr>
<tr>
<td>Bochum11673</td>
<td>45nm</td>
<td>5nm GaAs:C doped</td>
<td>(n = 4.5 \times 10^{11} \text{ cm}^{-2})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>C-delta doping</td>
<td>(\mu = 130'000 \text{ cm}^2/\text{Vs})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>15nm Al_{0.35}Ga_{0.65}As</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>25nm Al_{0.35}Ga_{0.65}As</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>650nm GaAs</td>
<td></td>
</tr>
<tr>
<td>Bochum11678</td>
<td>120nm</td>
<td>5nm GaAs:C doped</td>
<td>(n = 1.4 \times 10^{11} \text{ cm}^{-2})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>C-delta doping</td>
<td>(\mu = 180'000 \text{ cm}^2/\text{Vs})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>60nm Al_{0.35}Ga_{0.65}As</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>55nm Al_{0.35}Ga_{0.65}As</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>650nm GaAs</td>
<td></td>
</tr>
<tr>
<td>Bochum11680</td>
<td>135nm</td>
<td>5nm GaAs:C doped</td>
<td>(n = 1.8 \times 10^{11} \text{ cm}^{-2})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>C-delta doping</td>
<td>(\mu = 130'000 \text{ cm}^2/\text{Vs})</td>
</tr>
<tr>
<td></td>
<td></td>
<td>60nm Al_{0.35}Ga_{0.65}As</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>70nm Al_{0.35}Ga_{0.65}As</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>650nm GaAs</td>
<td></td>
</tr>
</tbody>
</table>

Bochum1282 and Bochum11673) [69]. For this purpose we have fabricated M-shaped Hall-bars (Fig. 3.2) with two of them aligned along the main crystallographic directions (angles \(\phi = 0^\circ\) and \(\phi = 90^\circ\)), and the other two rotated by \(45^\circ\) with respect to the main crystallographic directions (angles \(\phi = 45^\circ\) and \(\phi = -45^\circ\)). The angle \(\phi\) is defined in Fig. 3.2. The width of the Hall-bars is 25 \(\mu\)m, while the separation between neighboring contacts is 50 \(\mu\)m. We measured the longitudinal and transversal resistivity for each of these directions at a temperature of 4.2 K, and from there deduced the mobility anisotropy of the samples. The results are shown in the following table (\(\mu_{\text{high}}\) is the higher and \(\mu_{\text{low}}\) is the lower of the two mobilities along the main crystallographic directions, while the average mobility is given by \(\mu_{\text{average}} = (\mu_{\text{high}} + \mu_{\text{low}})/2\)).

The measured mobilities along the directions \(\phi = 45^\circ\) and \(\phi = -45^\circ\) are larger than \(\mu_{\text{low}}\) and smaller than \(\mu_{\text{high}}\), as expected. We see from the obtained results that the mobility anisotropy in all of the three investigated C-doped wafers is significantly smaller than the common mobility anisotropy in Si-doped (311) wafers [62]. Therefore C-doped p-type GaAs/AlGaAs heterostructures show superior electronic properties compared to Si-doped samples.
an angle $\phi$. The width of a Hall bar is 25 $\mu$m, while the separation between neighboring leads is 50 $\mu$m. The angle of the Hall-bar is defined in the figure.

**Table 3.2:** Scheme of the M-shaped Hall-bar used for measurements of mobility anisotropy. The width of the Hall-bar is 25 $\mu$m, while the separation between neighboring contacts is 50 $\mu$m. The angle of the Hall-bar is defined in the figure.

<table>
<thead>
<tr>
<th>wafer</th>
<th>direction ($\phi$)</th>
<th>mobility $\text{[cm}^2\text{/Vs]}$</th>
<th>$(\mu_{\text{high}} - \mu_{\text{low}})/\mu_{\text{average}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Bochum12029</td>
<td>$0^\circ$</td>
<td>106'600</td>
<td>0.013</td>
</tr>
<tr>
<td></td>
<td>$90^\circ$</td>
<td>108'000</td>
<td></td>
</tr>
<tr>
<td>Bochum1282</td>
<td>$0^\circ$</td>
<td>124'500</td>
<td>0.21</td>
</tr>
<tr>
<td></td>
<td>$90^\circ$</td>
<td>100'600</td>
<td></td>
</tr>
<tr>
<td>Bochum11673</td>
<td>$0^\circ$</td>
<td>93'500</td>
<td>0.37</td>
</tr>
<tr>
<td></td>
<td>$90^\circ$</td>
<td>134'300</td>
<td></td>
</tr>
</tbody>
</table>

### 3.2 Steps for fabricating semiconductor nanostructures

In order to make nanostructures accessible to external macroscopic measurement devices we need to bridge several length-scales. This path from the initial plain GaAs/AlGaAs heterostructure to the final device connected to the measurement setup (Fig 3.3(a)) involves several steps which will be explained in this paragraph. In the first step a square mesa structure of $20 \times 20 \mu m^2$ (Fig. 3.3(d)), together with the connecting leads (Fig. 3.3(b)) is defined with optical lithography. The second optical lithography step, followed by the evaporation of contact metallization and lift-off is used for defining contact pads with a diameter of 250 $\mu$m (yellow circles in Fig. 3.3(b)). After that AFM lithography is used for defining a nanostructure on the mesa (Fig. 3.3(f)). In addition, thick insulating oxide lines between the
3.3 Optical lithography

Optical lithography can be used for the fabrication of structures larger than $1 - 2 \mu m$. As mentioned in the previous paragraph, optical lithography is used in this thesis for the fabrication of Hall-bars and $20 \times 20 \mu m^2$ mesa structures, which serve as platforms for AFM lithography, as well as for defining the Ohmic contacts. The scheme of the standard photolithography process is given in Fig. 3.4, and the details about the

Figure 3.3: Bridging different length-scales: (a) Sample bounded into the chip carrier and placed in the sample-holder. (b), (c), (d) Scheme of the sample with mesa structure on different length-scales. (e) AFM micrograph of the mesa, after the complete structure is written, with clearly visible connecting oxide lines. (f) AFM micrograph of the nanostructure.
parameters used in the process can be found in Appendix B. In brief, after cleaning the wafer with aceton, isopropanol and water, the sample is covered with a positive photoresist and exposed to the UV light through a mask containing the desired pattern. The positioning of the mask is done with a commercial Karl Suss MJB3 mask aligner. After development, the photoresist exposed to UV light is removed, and the wafer outside the desired pattern is then etched down by 50-100 nm with H$_2$O:H$_2$SO$_4$:H$_2$O$_2$. In such a way, the 2DHG is removed in the etched regions. In the second photolithography step, used for the definition of the Ohmic contacts, the negative photoresist is spun on the sample surface and a different mask with contact patterns is aligned with respect to the predefined structure. After exposure and development of the photoresist, instead of etching, the contact metallization (Au/Zn/In) is evaporated onto the surface, and subsequently removed from the regions covered with photoresist by a lift-off process in acetone.

The resolution which can be achieved with optical lithography is limited by the wavelength of UV light and the optical alignment of the mask with respect to the sample. As mentioned this limit is around 1µm, and therefore it is clear that other techniques, like electron beam lithography or AFM oxidation lithography must be used for the fabrication of structures smaller than 1µm.
3.4 Ohmic contacts to two-dimensional hole gases

Low resistance Ohmic contacts to GaAs are crucial for performing low-noise transport measurements in nanostructures. While the technology for making good Ohmic contacts to n-type GaAs is well established, with Au/Ge/Ni serving as a contact metallization, the fabrication of Ohmic contacts on p-type GaAs appeared to be more difficult. Typically AuZn, AuBe or InZn metallizations are used for Ohmic contacts to p-type GaAs [70].

As mentioned above, we investigated five different p-type GaAs heterostructures. Our experience shows that making Ohmic contacts to deeper 2DHGs (heterostructures Bochum1282, Bochum11678 and Bochum11680) is straightforward, while making good Ohmic contacts to shallow 2DHGs (heterostructures Bochum12029 and Bochum11673) was quite troublesome and required a lot of optimizing.

Before each evaporation of Ohmic contact metallization the sample surface is cleaned for 2 minutes with plasma ashing (power 200 W) and the sample is then immediately placed in an evaporation machine. In the case of the deep 2DHGs we evaporate the following metallization directly on the sample surface: 40nm Au, 40nm Zn, 200nm Au and after evaporation we anneal the contacts at 480-500°C in 95% N₂ and 5% H₂ atmosphere for 2 minutes. For such contacts we measure typical 2-terminal contact resistances across the sample to be 200-300 kΩ at room temperature, and when the sample is cooled to 4.2K the contact resistance drops to around 10 kΩ and shows good linearity. Here it is important to emphasize that the quoted resistance includes two contacts and the part of the heterostructure between the contacts.

In the case of shallow 2DHG heterostructures Bochum12029 and Bochum11673 using only Au and Zn did not give satisfying results, and therefore we also added In in the contact metallization. We also found that in case of the thick initial layer of gold (40 nm) the contacts did not work and we had to reduce this thickness to 2 nm. Therefore we use the following contact metallization for shallow 2DHG samples: 2nm Au, 40nm Zn, 40nm In, 200nm Au annealed at 480-500°C in 95% N₂ and 5% H₂ atmosphere for 2 minutes. The 2-terminal resistances of these contacts are 300-500 kΩ at room temperature and they drop very fast upon cooling the sample to 4.2K to values around 10-20 kΩ. Here it is important to point out that the room temperature value of the 2-terminal contact resistance cannot give a clear answer if the contacts are good at low temperature. For some samples with similar values of the room temperature contact resistance, the contact resistance increased upon cooling to 4.2K. Therefore it is crucial to check if the contacts are good at 4.2K, before proceeding with the fabrication of nanostructures with AFM oxidation lithography.

We can mention here that in most reported transport experiments on bulk p-type GaAs heterostructures, the Ohmic contacts are made by soldering InZn. However, this approach is not suitable in our case, since the pieces of soldered InZn can stick out quite high and therefore prevent the AFM tip from approaching the sample sur-
face during the later processing of nanostructures with AFM oxidation lithography.

Finally, it is important to mention that indium becomes superconducting at $T_{\text{C,In}} = 3.4$ K, and zinc at $T_{\text{C,Zn}} = 0.85$ K. Therefore, for the measurements in a dilution fridge at the temperatures in the mK range the possible transition of the Au/Zn/In contact metallization into a superconducting state might be relevant for transport measurements, and this will be discussed in detail in chapter 4.

3.5 Metallic Schottky gates on p-type samples - hysteresis and sample instability

Semiconductor nanodevices are conventionally fabricated using electron beam lithography defined metallic split-gates, evaporated on top of the heterostructure. Applying voltages to these split-gates allows the control and manipulation of the confining potential in which the charge carriers reside. This approach of sample fabrication requires that upon evaporation of metal onto the semiconductor surface a good rectifying junction with a Schottky barrier forms. While split-gate devices fabricated on n-type GaAs/AlGaAs heterostructures generally show excellent electronic properties, split-gate devices on p-type GaAs/AlGaAs heterostructures show significant gate instabilities [21, 71]. Due to these technological difficulties, p-type nanostructures remained quite unexplored compared to their n-type counterparts. The reason why metallic gates are more leaky on p-type GaAs and why they induce more sample instabilities might be the fact that the Fermi level is pinned closer to the bottom of the valence band than to the top of conduction band and therefore the Schottky barriers are lower on p-type compared to n-type GaAs [72]. Also the Shottky barrier is operated in forward bias direction if the hole gas is depleted.

![Figure 3.5: I-V characteristics of the metallic top-gate on the wafer Bochum1282 with a 2DHG 100 nm below the surface at T=100 mK.](image)

We explored the behavior of metallic top-gates on five different p-type heterostructures quoted in the table in paragraph 3.1. In all cases the top-gates consist
of 10 nm of Ti evaporated directly on the sample surface, followed by 100 nm of Au. We find that metallic top-gates on the heterostructures with a deeper 2DHG (wafers Bochum1282, Bochum11678, Bochum11680) show good rectifying behavior at low temperatures (Fig. 3.5 shows an example of I-V curve of the top-gate on the wafer Bochum1282 at 100 mK), and can be used for tuning the hole density in these 2DHGs. On the other side metallic top-gates on the heterostructures with shallower 2DHGs (Bochum12029 and Bochum11673) are leaky even at low temperatures, showing almost completely Ohmic instead of rectifying properties and therefore cannot be used for tuning the density in these samples. In order to improve the performance of top-gates on these shallow wafers it would be adequate to investigate a possibility that instead of evaporating metals directly onto the sample surface, a thin insulating layer (for example SiO$_2$ or calixarene) is inserted between the sample surface and the metallic gate. This has not been done in this thesis and remains as a possibility for future work.

![Figure 3.6](image)

**Figure 3.6:** (a) Density relaxation in the heterostructure Bochum11680, after the top-gate is gradually changed from 0 to 0.1V. (b) Hysteresis in a QPC current upon sweeping the voltage applied to the QPC split-gates up and down. Inset: A micrograph of the split-gate device fabricated on the wafer Bochum11680: the results here refer to the left QPC, while the other gates are kept grounded during this measurement.

In the case of the three wafers with working top-gates (Bochum1282, Bochum11678, Bochum11680) we characterized the time until the density in a 2DHG reaches equilibrium after the value of top-gate voltage is changed. In the case of the wafer Bochum1282 this time is of the order of few seconds, while in the wafer Bochum11680 it is around 1000 s (Fig. 3.6(a)), and even longer in the wafer Bochum11678. This might not be a critical issue if one only wants only to change the global density in a 2DHG by setting a top-gate voltage at a fixed value. However, if one wants to tune a device by sweeping in time the voltages applied to the topgates, hysteretic effects appear. An example of such hysteretic behavior is observed in a measurement of a
current through a split-gate defined QPC upon sweeping the voltage applied to the
top-gates up and down (Fig. 3.6(b)). (In this measurement the voltage is applied to
the left pair of QPC gates, shown in inset of Fig. 3.6(b), while the remaining four
gates are kept grounded). A similar hysteretic behavior in response of the 2DHG to
the top-gate changes was observed in an extensive study of several different p-type
GaAs heterostructures performed in the Group of Prof. Wegscheider in Regensburg
[73]. However, the physical origins of such an expressed hysteresis in p-type GaAs
heterostructures are not yet clear.

Due to these difficulties observed in devices fabricated with metallic split-gates
we decided to employ another technique, namely Atomic Force Microscope (AFM)
oxidation lithography for the fabrication of p-type GaAs nanostructures.

3.6 AFM oxidation lithography

Atomic Force Microscope (AFM) oxidation lithography is a nanofabrication tech-
nique based on the possibility to locally oxidize surfaces of semiconductors or metals
with a moving metallic tip, acting as cathode. This technique is devised by Dagata
[74] and is later used for local oxidation of various substrates. It was first shown
in our group that an oxide line written on the surface of the semiconducting hetero-
structure completely depletes 2DEG located 50 nm or less below the surface [22].
A detail account about the fabrication of nanostructures with this technique can be
found in [75], and we give here a short overview.

All nanostructures in this thesis are written on the same wafer, Bochum12029,
with a 2DHG 45 nm below the surface. We use a commercial Asylum Research
MFP-3D AFM, placed in a closed chamber with a controlled atmosphere. The
scheme of the setup is shown in Fig. 3.7(a). Before starting AFM lithography,
the sample surface is cleaned with plasma etching for 2 min. This procedure burns
organic contamination on the surface and ensures good surface roughness (around
1nm peak-to-peak), which is important for AFM lithography. The sample is then
placed in a grounded sample holder below the AFM head, and the AFM chamber
is closed. The humidity in the chamber is adjusted to a value $43 - 45\%$ at room
temperature, by introducing in a controlled way moistened nitrogen through a water
bubbler into the AFM chamber. This ensures the formation of a thin water film
on the surface, which serves as an electrolyte for oxidation. Once the humidity
stabilizes, it is important to wait an additional 30-60 min before starting with AFM
lithography, to be sure that all thermal drifts are minimized.

During oxidation the AFM is operated in tapping mode with a cantilever being
excited to its resonance frequency of around 300 kHz. Nanostructures in this thesis
are written with Si cantilevers covered with a metallic Ti/Pt coating. However, it is
important to mention that the writing strongly depends on the microscopic details
of the tip, and that even nominally identical tips from the same batch show very
different writing properties.
The advantage of AFM oxidation lithography over other techniques is that once the sample is placed in the AFM chamber one can, with the same tip, inspect the sample surface, write the structure at a desired location and finally check the topography of the written structure. This method also allows later reparations of the structures.

Initially, AFM oxidation has been performed with a constant dc voltage being applied to the AFM tip [75]. Recent studies by Graf et al. [76] show that the application of a square ac-modulated voltage to the AFM tip leads to better reproducibility of the oxidation process and also enhances the aspect ratio of the oxide lines which improves their insulating behavior. Therefore the structures in this thesis are written with an ac-modulated voltage applied to the tip. Typical parameters used for AFM writing of the structures presented in this thesis are (for clear definition of the given parameters see [76]): humidity 43–45%, set-point 0.05–0.1, dc-offset voltage $V_{dc}$ in the range $-12.5\text{V}$ to $-22.5\text{V}$, square ac-modulation voltage $V_{ac}$ with a peak to peak value in the range 15V to 25V, frequency of the ac-modulation 750 – 1000 Hz, reset time 30 – 40% of the total period of ac modulation and the tip velocity 100–200 nm/s. It is important to point out that $V_{ac}$ was always chosen to be slightly larger than the absolute value of $V_{dc}$. Therefore, during the oxidation process the tip approaches the surface much more closely (set point 0.05-0.1) than during scanning (set point 0.7-0.8). In such a way oxide lines with a height of 15-18 nm and width around 100 nm could be written (Figs. 3.8(a) and 8.1(b)). Our tests show that for a given wafer with a 2DHG 45 nm below the surface, the minimum height of the

Figure 3.7: (a) Scheme of AFM oxidation lithography. (b) Explanation of the local depletion below the oxide line. (c) Transfer of the oxide pattern into the 2DHG.
oxide line which depletes the 2DHG beneath is around 12 nm.

It was shown in [75] that the growth of oxide lines is such that the oxide height above the sample surface is approximately equal to the depth of the oxide grown below the sample surface. We also show that in the case of a p-type GaAs wafer, the depth of the groove after removing the oxide with HCl, is approximately the same as the height of the oxide grown on top (Fig 3.8). It is observed that the electronic properties of AFM written structures do not change significantly after removing the oxide. This confirms that the oxide itself is irrelevant for the electronic properties of AFM written structures, and that AFM lithography represents an effective way to locally etch down the wafer to the depth approximately equal to the height of the oxide protruding the surface.

![Image](image_url)

Figure 3.8: (a) AFM defined QPC with a height of the oxide lines of around 15 nm. (b) QPC after removing the oxide with HCl: grooves with a depth similar to the height of the oxide lines remain.

A schematic explanation why etching down the surface of a wafer leads to the local depletion of the underlying 2DHG is given in Fig. 3.7(b). GaAs has the particular property that its Fermi level is pinned by the surface states to a fixed position above the edge of the valence band. When the wafer is slightly etched down, the sample surface is effectively brought closer to the 2DHG. Since the position of the valence band edge at the surface has to remain at the same position below the Fermi level, due to Fermi level pinning, the valence band is pushed down in energy (Fig. 3.7(b)), and the 2DHG becomes depleted below the oxide line. The insulating behavior of the oxide lines at 4.2K is demonstrated by measuring the I-V characteristic across the line (Fig. 8.1(c)) - the leakage current across the line was less than 1 pA for the applied voltages in the range [-500 mV, +400 mV]. However, these voltage ranges in which the oxide lines remain insulating may vary slightly from sample to sample.

Fig 3.7(c) shows how the pattern of the oxide lines at the surface of the wafer transfers into the pattern of insulating lines in the plane of a 2DHG. In such a way...
the desired structure is defined in the plane of a 2DHG. The regions of a 2DHG electrically separated by oxide lines, can serve as independent in-plane gates for tuning the structures [77].

3.7 Cryogenics

In order to be able to observe various quantum phenomena in semiconductor nanostructures, characteristic energy scales in the system have to be larger than the thermal energy $k_B T$. Therefore electronic transport experiments typically have to be performed at very low temperatures.

![Diagram of the 3He/4He dilution refrigerator.](image)

Figure 3.9: (taken from [34]) Scheme of the $^3$He/$^4$He dilution refrigerator.

Three types of refrigeration systems have been used in the experiments presented in this thesis. The first system is a simple dewar filled with liquid helium ($^4$He) under normal atmospheric pressure, which cools the sample to 4.2K. This system is equipped with a 5T magnet. The second system is also a $^4$He system with a
variable temperature insert (VTI) and in this system the temperature can reach values down to 1.6K. This is achieved by continuous pumping on the liquid helium and a reduction of the $^4$He vapor pressure down to values of around 10 mbars. This system is equipped with a 9T magnet. These two refrigeration systems are mostly used for the basic characterization of the samples.

Since holes in GaAs have effective masses $m^*$ several times larger than electrons, the typical energy scales in p-type samples are scaled down by a factor $1/m^*$ and are much smaller than characteristic energies in n-type systems. Therefore measuring at very low temperatures, in the mK temperature range, is essential for studying hole transport. Such low temperatures can be achieved in a dilution refrigerator. In our Oxford dilution refrigerators the temperature can reach values down to around 60 mK. This system is equipped with a 12T magnet.

The scheme of the $^3$He/$^4$He dilution refrigerator is shown in Fig. 3.9. The mixing chamber where the cooling down to 60 mK occurs is surrounded by an inner vacuum chamber, $^4$He main bath at 4.2K, a liquid nitrogen shield at 77K and an outer vacuum chamber. If the $^3$He/$^4$He mixture, with a $^3$He concentration larger than 6% is cooled below 0.87K, the condensed mixture will separate into two phases: the $^3$He poor phase (dilute phase) and a $^3$He rich phase which floats on the $^4$He rich phase. Cooling takes place at the phase boundary which is located in the mixing chamber. By pumping on the $^3$He poor phase in the still at a temperature below 1K practically only $^3$He is removed from the liquid, because the partial pressure of $^3$He is around 1000 larger than the partial pressure of $^4$He under these conditions. In order to keep equilibrium concentration of $^3$He in a $^3$He poor phase, $^3$He atoms are forced to flow across the phase boundary from $^3$He rich to $^3$He poor phase. This process provides cooling at the phase boundary due to the latent heat associated with this transfer. After being removed from the still, the $^3$He gas is again condensed in the 1K pot into the $^3$He rich phase and therefore $^3$He gas continuously circulates in the system. A detailed explanation of the dilution refrigerator working principle is given in [78].

### 3.8 Electrical setups for transport measurements

Depending on the characteristic resistance of the investigated sample, two measurement setups (Fig. 3.10) are typically used for transport measurements. For high resistance samples ($R \gg h/e^2$), two terminal measurements of the sample conductance are usually performed using a home-built I-V converter (Fig. 3.10(a)). A low dc or ac bias voltage (in the range of 1-100 µV) is applied symmetrically across the sample and the current through the sample is measured with an I-V converter with an amplification factor of $10^8$ or $10^9$ V/A. In order to suppress the noise from the cabling, low biases applied to the samples are obtained by scaling down the voltages from the ac or dc voltage sources with voltage dividers mounted directly at the top of the fridge. Besides, low-pass filters are included in the wiring of the sample.
3.8. Electrical setups for transport measurements

Figure 3.10: (taken from [34]) (a) Scheme of the two-terminal conductance measurement setup using an I-V converter. (b) Scheme of the four-terminal resistance measurement setup using a lock-in amplifier.

holders to prevent sample heating through the leads.

Low resistance samples \( R \leq \frac{\hbar}{e^2} \) are measured in a four-terminal setup (Fig. 3.10(b)). A low ac current (1-100 nA, depending on the particular measurement) with a low-frequency of 31 Hz is applied through the sample and the corresponding voltages are measured with a lock-in technique, using a Stanford SRS 830 lock-in amplifier. In very sensitive measurements additional low-noise voltage amplifiers, with a gain of 1000, were mounted directly at the top of the fridge, and the relevant signal was amplified before being measured with a lock-in amplifier. All measurements were automatized with LabVIEW computer programs.
Chapter 4

Magnetoresistance of In/Zn/Au Ohmic contacts

In this chapter we present experimental observations which can be related to the fact that indium, which is used for making Ohmic contacts to 2DHGs, becomes superconducting at mK temperatures at which our transport experiments are performed. As it was discussed in paragraph 3.4 we have used two types of Ohmic contacts to 2DHGs. For the heterostructure Bochum1282 with a 100nm deep hole gas good Ohmic contacts could be made using Zn and Au metallization, while for the shallow heterostructure Bochum12029 we had to use In, Zn and Au metallization to fabricate reliable Ohmic contacts. In both cases the contacts are annealed at 480-500°C in a 95% N\(_2\) and 5% H\(_2\) atmosphere for 2 minutes. Here it should be mentioned that indium becomes superconducting at \(T_{C,\text{In}} = 3.4\) K, and zinc at \(T_{C,\text{Zn}} = 0.85\) K. Therefore it is not unexpected that superconductivity effects appear at temperatures below 100 mK where the experiments are performed. However, since Au, which is not a superconductor is also present in contact metallization, we can expect that superconductivity will be partially suppressed due to the presence of Au, and that the critical temperature of the contacts alloy will be reduced compared to the critical temperatures quoted for clean materials. It is important to point out that our InZn contacts are not soldered, as it is commonly done on p-type samples, but they are evaporated as thin films in high vacuum, and therefore presumably preserve a high degree of cleanliness, which might be relevant for the appearance of superconductivity effects.

The observations which are discussed in this chapter are found in samples with In/Zn/Au contacts. Samples with Ohmic contacts containing only Zn/Au did not show any of the observed effects and displayed standard behavior. This clearly points to the presence of In as being responsible for the observed effects, which is in agreement with the fact that it is a better superconductor than Zn. The data presented in this chapter originate from measurements on one sample. Qualitatively the same behavior is observed on several other samples with In/Zn/Au contacts being evaporated in separate evaporation runs. However, the magnitudes of the
observed effects varied from one sample to another which suggests that these effects depend on microscopic details of the evaporated metallization.

It is important to emphasize that all presented effects are observed only in two-terminal and not in four-terminal resistance measurements which clearly indicates that the effects are related to the contacts and not to the 2DHG.

We have measured the two-terminal conductance between two-contact pads (yellow circles in Fig. 3.3(b)) on a $20 \times 20\mu m^2$ mesa. The current flows from the contact metallization, down through the part of the heterostructure where the contacts to the 2DHG form and then after passing the 2DHG goes in the same fashion out to the second contact. Therefore we basically deal with transport through the Superconductor/Semiconductor (bad conductor)/Metal (2DHG)/Semiconductor/Superconductor junction. A small ac bias of $10\mu V$ is applied between the contact pads and the current is measured with an I-V converter. We have reproduced the same observations for an applied dc voltage.

The following distinct features are observed (Fig. 4.1-4.4):

- a strong reduction of the current (a dip in conductivity) at $B=0T$ at fixed bias voltage;
- hysteresis in magnetic field sweeps up and down
- non-linearity of I-V characteristic of the contacts, which manifests as a dip in a differential conductance at zero bias.

All tested pairs of contacts show the same behavior.

Fig. 4.1 shows the measured current across a pair of Ohmic contacts as a function of magnetic field at different temperatures indicated in the figure. The B-sweep rate is $0.1T/min$, and a fixed ac bias of $10\mu V$ was applied. The B-field was continuously swept up and down in the range (-0.25T, 0.25T), and the sweep direction was immediately changed once the borders of the given interval are reached. For each temperature 6 sweeps were performed: the 3 traces from the up sweeps (red traces) lay nicely on top of each other, as well as the 3 traces from the down sweeps (blue traces). However, in each of the two groups of traces there is an expressed asymmetry between the part below $B=0T$, and above $B=0T$ (Fig. 4.1). It is important to point that the part below $B=0T$ for sweep up traces is symmetric with respect to the part above $B=0T$ for sweep down traces and vice versa, the part above $B=0T$ for sweep up traces is symmetric to the part below $B=0T$ for sweep down traces (Fig. 4.1). Therefore, a clear hysteretic behavior is present in the current through the Ohmic contacts upon sweeping the magnetic field up and down. One can see that for both sweep directions when the sweep starts, the current changes along the lower of the two branches (Fig. 4.1) and once the point $B=0T$ is crossed it jumps to the upper branch. When the sweep stops at $B=0.25T$ (or $B=-0.25T$) the current instantaneously drops from the upper to the lower branch. Such a behavior suggests that the observed asymmetry between the lower and the upper branch on a single trace is caused by the fact that upon approaching the point $B=0T$, the contacts experience a transition to the superconducting state, while after crossing $B=0T$ they experience a transition out of the superconducting state.
Chapter 4. Magnetoresistance of In/Zn/Au Ohmic contacts

Figure 4.1: The measured current across a pair of Ohmic contacts as a function of magnetic field. Red traces correspond to B-field sweeps up and blue traces to sweeps down — hysteretic behavior is present. The sweep rate is 0.1 T/min, and an ac bias of 10 µV is applied across the contacts. The measurements are performed at different temperatures indicated on the right side. Note that the vertical offsets between the traces for different temperatures are not artificially introduced, but originate from the change of contact resistance with temperature.

As the temperature is increased, the hysteresis becomes weaker and eventually disappears above 300 mK (Fig. 4.1). This indicates that hysteresis is related to the superconducting contacts and that it disappears as the superconductivity of the contacts is destroyed. It should be pointed out that the offsets between the traces for different temperatures in Fig. 4.1 are real and originate from the change of contact resistance with temperature which will be discussed later.

We also observe a strong decrease in current through the contacts (dip in conductance) at B=0T. The typical width of the peak is around 20 mT at T=65 mK. We confirm that this conductance dip originates from the contacts and not from a 2DHG, by measuring 4-terminal resistance of the sample. The four terminal resistance reveals weak antilocalization, which means an increase of the conductance at B=0T (see chapter 6). The observation of the conductance peak with a width less than 0.5 mT in a 4-terminal measurement (chapter 6) in contrast to the conductance dip with a width around 20 mT in a 2-terminal measurement demonstrates that the conductance dip in a 2-terminal measurement indeed originates from the contacts and not from a 2DHG. In Fig. 4.1 we can see that the conductance dip at B=0T becomes weaker and narrower as the temperature is increased and completely disappears above T=230 mK. Such a behavior suggests that the conductance dip is related to superconductivity of the contacts. The magnetic field of around 10
mT below which the conductance experiences a sudden drop at T=65 mK might correspond to the lower critical field $H_{c1}$ of the In/Zn/Au contact alloy.

![Graph](image)

**Figure 4.2:** Current through the pair of Ohmic contacts at T=65mK as a function of magnetic field at fixed ac bias of 10µV and B-sweep rate 0.1T/min. Red trace correspond to B-field sweeps up and blue traces to sweep down. As the B-field approaches the values around 400mT, the hysteresis disappears.

We have further explored disappearance of the hysteresis at higher magnetic fields. In Fig. 4.2 we present the current through the contacts as the magnetic field is swept back and forth in the range (-0.25T, 0.5T). It can be seen that after going beyond the field B=0.25T, the hysteresis gradually decreases and eventually completely disappears at the fields above B=0.4T. This field might be related to the upper critical field $H_{c2}$ at which the In/Zn/Au alloy experience the transition from the so-called mixed state to the normal state.

The dependence of the hysteresis on the magnetic field sweep rate $dB/dt$, at a fixed temperature of 65 mK, is investigated in Fig. 4.3. We observe that upon decreasing the sweep-rate from 0.1 T/min the hysteresis weakens and for the sweep-rate of 0.01T/min becomes almost completely suppressed.

Finally, we have measured the I-V characteristics of the Ohmic contacts at different temperatures (Fig. 4.4(a)). At T=65 mK a nonlinearity around zero bias is observed, which becomes weaker as the temperature increases and almost completely disappears above 300 mK. In order to visualize these data in a more plausible way we plot in Fig. 4.4(b) the differential conductance $dI/dV$, obtained by numerical differentiation of the measured I-V curves, as a function of bias. It can be seen that at T=65 mK a very expressed conduction dip develops around zero bias with a minimum around $1 \times 10^{-5} \Omega$, and this dip weakens fast as the temperature increases up to 340 mK. If we analyze this change quantitatively we see that at T=340 mK, the low-bias differential resistance is around 30 kΩ and it increases to around 100
Chapter 4. Magnetoresponse of In/Zn/Au Ohmic contacts

0.1 T/min
0.05 T/min
0.025 T/min
0.01 T/min

Figure 4.3: Current through a pair of Ohmic contacts at $T=65\,\text{mK}$ as a function of magnetic field. A fixed ac bias of 10$\mu\text{V}$ is applied. The red traces correspond to B-field up-sweeps and blue traces to down-sweeps. The measurements are performed at different B-field sweep rates indicated in the figure. Traces are offset vertically for clarity.

k$\Omega$ as the temperature is reduced to 65 mK. Such a strong increase of the low-bias two terminal contact resistance as the temperature is lowered below 100 mK can be very harmful for low-noise transport measurements, which require relatively low contact resistances. However, we mentioned earlier that the magnitude of these superconductivity-related effects depends strongly on the microscopic details of the evaporated In/Zn/Au metallization, and that it might change from one evaporation to the other. In the samples where the superconductivity effects are not so strongly expressed, this temperature change of the low-bias contact resistance at very low-temperatures might not be so drastic, and low-noise transport measurements can be performed.

Let us mention that similar observations of the zero-bias dip in the differential conductance, as well as the strong reduction of the differential conductance as the magnetic field is reduced below 40 mT (conductance dip around $B=0\,\text{T}$) are reported by Marsh et al. \[79\] in the case of 2DEG contacted with pure indium contacts. These observations are found in the so-called clean-dirty transition in which the ratio between the 2DEG mean free path $l_\text{e}$, and the proximity-effect-induced pair coherence length $\xi_n$ in the normal material was around 1 ($l_\text{e}/\xi_n = 1.2$). However, the observation of the hysteresis in the current upon sweeping magnetic field is not reported in this paper.

To conclude, we have observed three distinct features in transport through In/Zn/Au Ohmic contacts fabricated on 2DHG: hysteresis in the 2-terminal current through the contacts upon sweeping magnetic field up and down, strong current
Figure 4.4: (a) Temperature dependence of the I-V characteristics of In/Zn/Au Ohmic contacts. (b) Temperature dependence of the differential conductance $dI/dV$, obtained by numerical differentiation of the I-V traces shown in (a).

suppression for the B-fields smaller than 10 mT (conductance dip), and non-linear I-V characteristic around zero bias. The observed temperature dependence of these three features suggests that they could be related to the superconductivity of the In/Zn/Au contact metallization, which is caused by presence of In. From these temperature and magnetic field dependencies of the current through the contacts we estimate the critical temperature of the In/Zn/Au alloy to be around $T_c \approx 300$ mK, and the critical magnetic fields $H_{c1} \approx 10$ mT and $H_{c2} \approx 400$ mT, which seems realistic.
Chapter 5

Magnetotransport in p-type carbon doped GaAs heterostructures

5.1 Introduction - Spin splitting of hole subbands due to inversion asymmetry

Energy bands in a periodic solid-state system possess the so-called Kramers degeneracy $E_{k\uparrow} = E_{-k\downarrow}$ due to the time-reversal symmetry ($k$ denotes the wave-vector and $\uparrow, \downarrow$ stands for the spin direction of the moving carrier). If in addition the system is characterized by spatial inversion symmetry, the relation $E_{k\uparrow} = E_{k\downarrow}$ holds and the two spin states are degenerate for any wave-vector $k$. Lack of inversion symmetry in a solid leads to a lifting of the spin degeneracy even in the absence of external magnetic fields. This lack of inversion asymmetry may be due to bulk inversion asymmetry (Dresselhaus splitting [44]) or due to structural inversion asymmetry (Rashba splitting [45]). Particularly the later mechanism received considerable interest recently due to the fact that the advances in material growth technologies give the opportunity to tailor the structural asymmetry of semiconductor heterostructures and therefore to tune the spin-splitting in these systems. This property makes such materials promising for development of spintronics.

Spin-splitting in heterostructures with inversion asymmetry can be attributed to the effective in-plane magnetic field $B_{\text{eff}}$, induced by strong spin-orbit coupling [43], which is present in the reference frame of the moving carriers. Although external magnetic fields are absent, an electric field, built in the heterostructure perpendicularly to the plane of a 2D electron (hole) system, is seen as a momentum dependent effective magnetic field $B_{\text{eff}}$ in the reference frame of the moving carriers. However, we should keep in mind that in non-magnetic material at zero external magnetic field, the net spin polarization has to be zero. Therefore, it is important to point out here that $B_{\text{eff}}$ varies in magnitude and orientation as a function of in-plane wave...
5.1. Introduction - Spin splitting of hole subbands due to inversion asymmetry

vector $\mathbf{k}$. Even though for a given $\mathbf{k}$ the states are spin polarized, after summing over all occupied $\mathbf{k}$ states the total spin polarization in each of the two spin-split subbands vanishes. In the following we will denote the two spin-split subbands by spin up and spin down, having in mind the spin orientation of the states with respect to the direction of the local effective magnetic field. Due to the different dispersion relations of the two spin-split subbands, their densities will be different and this difference $\Delta N$ will be called carrier imbalance in the further text.

![Figure 5.1: Self-consistently calculated dispersions $E_{\pm}(k_{\parallel})$ of the topmost HH subband for a (001) grown GaAs/AlGaAs heterostructure with a density $N = 2 \times 10^{11} \text{ cm}^{-2}$ (lower figure) together with the spin-splitting $E_{+}(k_{\parallel}) - E_{-}(k_{\parallel})$ (upper figure). $k_{\parallel}$ is along (100) direction in the plane of a 2DHG.](image)

As discussed in paragraph 2.2.3 the band structure of holes in 2DHGs is such that for typical densities ($10^{11} \text{ cm}^{-2}$) in these systems the Fermi level is placed above the top of the light hole (LH) subband (Fig. 2.7(a)), and only the holes from the heavy hole subband (HH), with a $z$-component of the total angular momentum $J_z = \pm 3/2$, contribute to the transport. In the presence of strong spin-orbit interactions this heavy hole subband is further split into two subbands with opposite $J_z$. Therefore a 2DHG with strong SO interaction effectively behaves as a two-subband system, with holes from spin-split heavy-hole subbands having opposite spins $+3/2$ or $-3/2$.

An example of self-consistently calculated dispersions $E_{\pm}(k_{\parallel})$ of the topmost HH subband for a (001) grown GaAs/AlGaAs heterostructure with a density $N = 2 \times 10^{11} \text{ cm}^{-2}$ together with the spin-splitting $E_{+}(k_{\parallel}) - E_{-}(k_{\parallel})$ is shown in Fig. 5.1 [43]. It can be seen that in this simulated structure the spin-splitting at the Fermi
level is around 0.5 meV and the Fermi energy is around 2 meV. It is also important to point out the non-linear spin-splitting of the heavy hole subband - it was shown by Gerchikov and Subashiev [49] that spin splitting of the heavy hole subband due to structural inversion asymmetry is cubic in the $k_\parallel$ vector ($\Delta E_{SIA}^{HH} \sim \pm \beta k_\parallel^3$), in contrast to the case of the conduction band where this splitting is linear in the $k_\parallel$ vector.

In transport experiments the existence of the spin-split subbands, induced by inversion asymmetry, results in a beating of the Shubnikov-de Haas (SdH) oscillations. The Landau levels of the two non-equally populated subbands give rise to magnetoresistance oscillations with slightly different frequencies, which are then superimposed for the total magnetoresistance. This effect was observed for the first time by Störmer et al. in a Be-doped (100) p-type GaAs/AlGaAs heterostructure [54]. Besides, spin splitting due to inversion asymmetry can also be characterized by weak antilocalization measurements, Raman scattering, spin resonance or cyclotron resonance measurements. Although the spin-splitting of the conduction band in GaAs/AlGaAs has been detected by weak antilocalozation [80] and Raman scattering [81], a beating of SdH oscillations in n-type GaAs/AlGaAs heterostructures has not been observed, except in a special sample, treated with e-beam irradiation [82]. This fact that beating in SdH oscillations is commonly observed for p-like valence band, but not for s-like conduction band of GaAs/AlGaAs heterostructures, indicates that SO interaction effects are stronger for holes than electrons. However, a beating in SdH oscillations was observed for electrons in other narrow-gap semiconductors like InGaAs, InAs or HgTe [83].

The possibility to tune the splitting of the two spin-split subbands with an external metallic top-gate was simultaneously demonstrated for electrons in InGaAs by Engels et al. [47] and Nitta et al. [48]. The effect of the electric field induced by a voltage applied to the top-gate is twofold: it changes the density of carriers in the quantum well and also changes the asymmetry of the quantum well, which is relevant for the inversion asymmetry induced spin-splitting. The tunability of the spin splitting for holes in (311) p-type GaAs/AlGaAs heterostructures was demonstrated by Lu et al. [84]. A detailed review of the relevant theories and experiments related to the beating of SdH oscillations due to inversion asymmetry induced splitting of subbands is given in [83].

In this chapter we report magnetotransport measurements on a novel type of (100) p-type GaAs/AlGaAs heterostructures obtained by doping with carbon. The high-quality of the two investigated heterostructures (Bochum1282 and Bochum12029) was demonstrated by the observation of highly resolved SdH oscillations, pronounced integer quantum Hall plateaus as well as characteristic features of fractional quantum Hall effect at filling factors 4/3 and 5/3 at temperatures around 70 mK. The observed beating pattern of low-field Shubnikov-de Haas oscillations in both samples represents clear evidence for the existence of two spin-split subbands which arise from strong spin-orbit interactions. By analyzing densities of the two spin-split subbands we deduce that the spin-orbit induced splitting of the heavy hole subband at
5.2. Basic characterization of p-type heterostructures

the Fermi level is around 30% of the Fermi energy in both heterostructures. We also analyze a classical positive magnetoresistance, together with a non-linear behavior in the Hall effect around \( B = 0 \) T, originating from the presence of the two spin-split subbands in the system. Finally we estimate the effective mass of the holes from the temperature dependence of the SdH oscillations.

5.2 Basic characterization of p-type heterostructures

We have performed Shubnikov-de Haas (SdH) and Hall measurements in two Hall-bars fabricated from the wafers Bochum1282 (100nm deep 2DHG) and Bochum12029 (45nm deep 2DHG). The Hall-bars are \( W = 100 \) µm wide, and the separation between the contacts used for the measurements of the longitudinal resistance is either \( L = 200 \) µm or \( L = 500 \) µm. In all plots the resistivity \( \rho_{xx} \), obtained after scaling the measured resistance by a geometrical factor \( W/L \), is presented.

As it was mentioned in chapter 2, three refrigeration systems are used throughout this thesis for sample characterization and examples of SdH and Hall measurements on the same sample (Bochum1282) in each of these systems at temperatures of 4.2 K, 1.7 K and 100 mK, respectively are shown in Fig. 5.2. At 4.2 K there are still no indications of SdH oscillations and only as the temperature is lowered to 1.7 K SdH oscillations start to develop with minima at filling factors 2 and 3. SdH oscillations and the plateaus in the quantum Hall effect are fully developed only in measurements in a dilution fridge at temperatures below 100 mK.

![Figure 5.2: Shubnikov de Haas and Hall measurements in the wafer Bochum1282 in three different refrigeration systems at the temperatures: (a) 4.2K, (b) 1.7K and (c) 100mK](image)

A Hall-bar fabricated from the wafer Bochum1282 with a 100 nm deep 2DHG, was equipped with a homogeneous Ti/Au top-gate, with the low-temperature I-V characteristic shown in Fig. 3.5. With this top-gate, the density in the sample could be tuned in the range from \( 1.7 \times 10^{11} \) cm\(^{-2}\) to \( 3 \times 10^{11} \) cm\(^{-2}\), while the corresponding mobility at 70 mK was in the range from 70’000 cm\(^2\)/Vs to 160’000 cm\(^2\)/Vs (Fig. 5.3(a)). The average mobility is determined from the resistivity at
zero magnetic field, $\rho_{xx}$. As it can be seen from Fig. 5.3(b) this zero-field resistivity keeps decreasing even below 1 K, and saturates only below 300 mK. Therefore, the sample mobility keeps increasing even at very low temperatures, and saturates below 300 mK. At the fixed density of $3 \times 10^{11}$ cm$^{-2}$ the mobility at 4.2K is 130’000 cm$^2$/Vs and it grows to 160’000 cm$^2$/Vs at 70 mK.

A similar behavior is observed in measurements on a Hall-bar fabricated from the wafer Bochum12029 with a 2DHG 45nm bellow the surface. There was no top-gate evaporated on this sample (see paragraph 3.5) and all the measurements have been performed at the fixed density of $3.8 \times 10^{11}$ cm$^{-2}$. The sample mobility at 4.2 K is 120’000 cm$^2$/Vs and it increases to 200’000 cm$^2$/Vs at 70 mK. Such a strong temperature dependance of holes’ transport characteristics, even for temperatures in mK range, originates presumably from the large effective mass of holes and the fact that the relevant energy scales are scaled down by a factor $1/m^*$ compared to electron systems. It is also reported in [85] that the strong temperature dependence of the holes’ resistivity is related to the spin-orbit induced splitting of the heavy hole subband - the larger the spin splitting, the stronger the temperature dependence of the resistivity is.

### 5.3 Integer and Fractional Quantum Hall effect

We have measured the longitudinal and Hall resistivity in both samples in a dilution fridge with a base temperature around 70 mK. Measurements are performed in a 4-terminal geometry with a small ac current of 10-20 nA applied across the sample. Figure 5.4 shows the longitudinal resistivity $\rho_{xx}$ (red traces) and the Hall resistivity $\rho_{xy}$ (blue traces) for both heterostructures (Bochum1282 and Bochum12029). The high quality of these carbon doped heterostructures is documented by the observation of pronounced integer quantum Hall plateaus, fractional quantum Hall states
4/3 and 5/3, as well as highly resolved low-field Shubnikov de Haas oscillations.

Figure 5.4: Longitudinal (red) and Hall (blue) resistivity at T=70 mK in two different carbon doped p-type GaAs/AlGaAs heterostructures: Bochum1282 with 100nm deep 2DHG (upper plot) and Bochum 12029 with 45nm deep 2DHG (lower plot). Insets show highly resolved low-field Shubnikov de Haas oscillations.

It is important to point out that in the Bochum1282 sample, beside the fractional states 5/3 and 4/3 which are observable both as minima in SdH oscillations and plateau-like structures in the Hall resistivity, the signature of a fractional state 5/2 is present as a minimum in SdH oscillations. The investigation of this fractional quantum Hall state came into focus recently, due to theoretical proposals for using this state for topologically protected quantum computing.

In the sample Bochum12029 the fractional states 5/3 and 4/3 manifest as minima in SdH oscillations, and as a dip-like and not plateau-like structures in the Hall resistivity. The same features are observed in another Hall-bar made from the same
Chapter 5. Magnetotransport in p-type carbon doped GaAs heterostructures

The inserts in Fig. 5.4 show highly resolved SdH oscillations at low B-fields in both heterostructures, which is another evidence for the high quality of the samples. They clearly show that at very low B-fields only one frequency is present in SdH oscillations, and as the B-field increases, the contribution from the second frequency becomes visible. These two frequencies originate from the two spin-split subbands and will be discussed in the next paragraph.

The fabrication of high mobility C-doped p-type GaAs heterostructures progress rapidly [66, 67] with the highest mobility reported so far being $2.2 \times 10^6$ cm$^2$/Vs [68]. The fact that carrier-carrier Coulomb interactions are much stronger in these systems relative to the Fermi energy, compared to n-type GaAs systems, makes C-doped 2DHGs ideal systems for studies of the so-called metal-insulator transition [85–88] and new phases in the quantum Hall regime [89, 90].

5.4 Beating of the Shubnikov-de Haas oscillations

We now focus on the analysis of the low-field magnetoresistance. As it was mentioned above, the Fourier analysis of the resistivity $\rho_{xx}$ vs. $1/B$ is often used to deduce the densities of the two spin-split subbands, originating from the inversion asymmetry of the heterostructures. The densities of the two spin-split subbands $N_{1,2}$ are related to the two frequencies $f_{1,2}$ obtained from the Fourier transform of SdH oscillations as $N_{1,2} = (e/h) \cdot f_{1,2}$ [48, 54, 84, 91].

Three B-field regions, where SdH oscillations exhibit a different behavior, can be identified in each of the two samples. As it can be seen in Fig. 5.5 and the insets of Fig. 5.4 in the regions of very low fields ((0.15T, 0.35T) for the sample Bochum12029 and (0.2T, 0.5T) for the sample Bochum1282) only SdH oscillations originating from one of the spin-split subbands, with a higher mobility, develop and there is only one peak in the Fourier spectrum related to the density of this subband.

As the B-field is further increased, the contribution from the second spin-split subband becomes visible, and in this intermediate B-field regime (up to approximately 2 T) three peaks are visible in the Fourier transform, corresponding respectively to the populations of each of the two spin-split subbands and the total density. Examples of Fourier transforms in this B-field regime are shown in Fig. 5.5. Fig. 5.5(b) shows the Fourier transform of the magnetoresistance of the sample Bochum1282, taken in the B-field range (0.4 T, 1.5 T), with the three peaks developing at 4.45 T, 7.7 T, and 12.3 T. Fig. 5.5(d) shows the Fourier transform of the magnetoresistance of the sample Bochum12029, taken in the B-field range (0.35 T, 1 T), with the three peaks developing at 5.4 T, 10.6 T, and 16 T. In both cases the relation $f_1 + f_2 = f_{tot}$, which reflects the fact that the two subband densities sum up to the total density, is reasonably satisfied.

For the B-fields above approximately 2 T we see the SdH oscillations which are related to the total density in the system, and only one peak, which determines the
5.4. Beating of the Shubnikov-de Haas oscillations

total density, is present in the Fourier transform.

In order to increase the resolution of Fourier transforms, we were using a standard procedure of padding data with zeroes, which involved subtraction of the smooth background, multiplication of the data with the Hamming window and adding zeroes at both ends of the data set. It should be mentioned that the peak positions were slightly dependent on the B-field range in which the Fourier transform is taken. These shifts were smaller than 5% of the obtained frequencies, so we also assume that the corresponding densities of the spin-split subbands are determined within 5% accuracy.

![Figure 5.5](image)

Figure 5.5: (a) Low field magnetoresistance of the sample Bochum1282, with a top gate grounded, so that the total density is around $3.0 \times 10^{11}$ cm$^{-2}$. (b) Fourier transform of the magnetoresistance shown in (a) taken in the B-field range (0.4T, 1.5T). (c) Low field magnetoresistance of the sample Bochum12029 with a density of around $3.8 \times 10^{11}$ cm$^{-2}$. (d) Fourier transform of the magnetoresistance shown in (c) taken in the B-field range (0.35T, 1T). Both Fourier transforms display 3 peaks, corresponding to the densities of the two spin-split subbands and the total density.

We further quantify the strength of the spin-orbit interactions in each of the two samples from the charge imbalance between the spin-split subbands. A cubic in
Chapter 5. Magnetotransport in p-type carbon doped GaAs heterostructures

$k$ wave vector dependence for the splitting of the heavy-hole subband in the (100) plane is assumed, $\Delta_{SO} = 2|\beta k|^3$, [43, 49], and the spin-orbit coupling parameter $\beta$ is calculated using the expressions 6.38 and 6.39 from the reference [43]. As it can be seen in Fig. 5.1 the wavevectors $k_1$ and $k_2$ at which the two spin-split subbands intersect the Fermi level are different - this difference becomes larger as the spin-orbit interaction is stronger. The energy splitting of the two spin-split subbands therefore depends on the $k$-vector where this splitting is calculated. The values of the spin-splitting energy which we quote further in the text, are all obtained using the smaller of the two wavevectors, and therefore represent the lower bound for the spin-splitting of the heavy hole subband.

![Figure 5.6](image)

Figure 5.6: The plots are related to the sample Bochum1282, in which the density can be tuned with a top-gate. (a) Charge imbalance in the spin-split subbands as a function of the total density. (b) Relative charge imbalance as a function of the density. (c) Spin-orbit splitting energy of the heavy hole subband as a function of the Fermi energy in the system and the total density. (d) Relative strength of the SO splitting energy, $\Delta_{SO}/E_F$ as a function of the Fermi energy.

For the sample Bochum12029 with a 2DHG 45nm below the sample surface and the total density $N = 3.8 \times 10^{11}$ cm$^{-2}$, the densities of the two spin-split subbands deduced from the Fourier peaks in Fig. 5.5(d) are $N_1 = 1.35 \times 10^{11}$ cm$^{-2}$ and $N_2 = 2.45 \times 10^{11}$ cm$^{-2}$, and the corresponding wavevectors where the two bands intersect the Fermi level are $k_1 = 1.3 \times 10^8$ m$^{-1}$ and $k_2 = 1.75 \times 10^8$ m$^{-1}$. The relative charge imbalance is $\Delta N/N = 0.29$. The spin-orbit coupling parameter is estimated to be
5.5. Classical positive magnetoresistance in two-subband systems

\[ \beta = 2 \times 10^{-28} \text{eVm}^3, \] which gives the spin-orbit induced splitting of the heavy hole subband at the Fermi level to be \( \Delta_{SO} \approx 0.85 \text{meV} \). Due to the large effective mass of holes the Fermi energy in the system is quite small, \( E_F = 2.5 \text{meV} \), and therefore the strength of the spin-orbit interactions relative to the kinetic energy is quite big \( \Delta_{SO}/E_F \approx 33\% \).

The second sample, Bochum1282 with a 2DHG 100nm below the surface, is equipped with a metallic top-gate and therefore the total density in the system could be tuned. When the top-gate is kept grounded the density in the system is \( N = 3 \times 10^{11} \text{cm}^{-2} \), and the densities of the two spin-split subbands are \( N_1 = 1 \times 10^{11} \text{cm}^{-2} \) and \( N_2 = 2 \times 10^{11} \text{cm}^{-2} \) giving the relative charge imbalance of \( \Delta N/N = 0.33 \). The corresponding wavevectors where the two bands intersect the Fermi level are \( k_1 = 1.1 \times 10^8 \text{m}^{-1} \) and \( k_2 = 1.6 \times 10^8 \text{m}^{-1} \). The spin-orbit coupling parameter is \( \beta = 2.5 \times 10^{-28} \text{eVm}^3 \), giving the spin-orbit induced splitting of the heavy hole subband of \( \Delta_{SO} \approx 0.7 \text{meV} \). The Fermi energy for a given total density is \( E_F = 2.0 \text{meV} \), so the strength of the spin-orbit interactions relative to the kinetic energy is \( \Delta_{SO}/E_F \approx 35\% \). The evolution of charge imbalance \( \Delta N, \) the relative charge imbalance \( \Delta N/N, \) the spin-splitting energy \( \Delta_{SO} \) at the Fermi level and \( \Delta_{SO}/E_F \) upon changing the total density in the system with a top-gate is shown in Fig. 5.6. We see that for all accessible densities the ratio \( \Delta_{SO}/E_F \) is quite large, in the range from 0.29 to 0.35, documenting the presence of strong SO interactions in C-doped GaAs/AlGaAs heterostructure. These measurements also demonstrate the possibility to tune the spin-orbit interaction in this sample with a metallic top-gate.

### 5.5 Classical positive magnetoresistance in two-subband systems

The longitudinal magnetoresistance of a system with two types of charge carriers with different mobilities, shows a parabolic behavior around zero magnetic field, while the Hall resistivity, beside the usual linear in B term contains a small, cubic in B correction at low fields. This is a purely classical effect and follows from standard Drude theory of conductivity [92]. If inter-subband scattering between the two subbands with the two different types of carriers is present, a more complex theory based on the Boltzmann transport equation has to be considered [93]. However, the qualitative behavior of the low-field magnetoresistance remains very similar to that obtained using the Drude model.

In a transport theory of two-subband systems with inter-subband scattering included, developed by Zaremba [93], the longitudinal and transversal magnetoresistivity are given as:

\[
\rho_{xx} = \frac{m^*}{e^2} \cdot \text{Re}\left(\frac{1}{\text{Tr}N(K - i\omega_c I)^{-1}}\right), \quad \rho_{xy} = \frac{m^*}{e^2} \cdot \text{Im}\left(\frac{1}{\text{Tr}N(K - i\omega_c I)^{-1}}\right) \tag{5.1}
\]

where \( \text{Tr} \) stands for the trace operation, \( \mathbf{I} \) is the \( 2 \times 2 \) unit matrix, \( \mathbf{N} \) is a matrix.
defined as $N_{ij} = \sqrt{N_i N_j}$ ($N_1, N_2$ are the densities of the two subbands), and $K$ is the scattering matrix defined as

$$
\begin{pmatrix}
K_1 & -K_{12} \\
-K_{12} & K_2
\end{pmatrix}
$$

where $K_1, K_2$ are the scattering rates within each of the two subbands, while $K_{12}$ is the inter-subband scattering rate between the two subbands.

Previously, a strong positive magnetoresistance was observed in p-type (311) GaAs heterostructures [94]. However, in that case the low-field magnetoresistance could not be fitted satisfactorily with the two-subband theory, even when intersubband scattering was taken into account, and this was attributed to the strong mobility anisotropy in (311) samples.

Here we study the classical positive magnetoresistance in two p-type (100) GaAs heterostructures. Figure 5.7 shows a strong positive magnetoresistance around $B=0$ T in the sample Bochum1282 in two configurations: (a1) $V_{TG} = 0$ V, $N=3.0 \times 10^{11}$ cm$^{-2}$ and (b1) $V_{TG} = 1$ V, $N=2.3 \times 10^{11}$ cm$^{-2}$ as well as in the sample Bochum12029 with a density $N=3.8 \times 10^{11}$ cm$^{-2}$ (c1). Full blue lines correspond to the measured data, while dashed red lines show the fits using the two-subband theory with intersubband scattering included (eq. 5.1). In the fitting procedure, the densities of the two subbands $N_1, N_2$ are taken from a Fourier analysis of the SdH oscillations as input parameters, while the scattering rates $K_1, K_2, K_{12}$ are fitting parameters.

For the sample Bochum1282 we obtain excellent fitting. In the configuration $V_{TG} = 0$ V, $N=3.0 \times 10^{11}$ cm$^{-2}$ (Fig. 5.7(a1)) the scattering rates of the individual subbands are $K_1 = 0.18 \times 10^{11}$ s$^{-1}$, $K_2 = 0.45 \times 10^{11}$ s$^{-1}$, while the inter-subband scattering rate is $K_{12} = 0.012 \times 10^{11}$ s$^{-1}$. We see that the intersubband scattering rate is much smaller compared to the scattering rates of the individual subbands. The corresponding mobilities of the two subbands are approximately (without taking into account $K_{12}$) $\mu_1=270'000$ cm$^2$/Vs and $\mu_2=110'000$ cm$^2$/Vs, which explains why in SdH measurements we first resolve only oscillations arising from the subband with population $N_1$, and oscillations from the subband with population $N_2$ become visible only at higher B-fields. In the second configuration $V_{TG} = 1$ V, $N=2.3 \times 10^{11}$ cm$^{-2}$ (Fig. 5.7(b1)) the scattering rates of the individual subbands are $K_1 = 0.33 \times 10^{11}$ s$^{-1}$, $K_2 = 0.64 \times 10^{11}$ s$^{-1}$, while the inter-subband scattering rate is $K_{12} = 0.086 \times 10^{11}$ s$^{-1}$. We see that the inter-subband scattering rate is still much smaller compared to the scattering rates of individual subbands. However, as the density is reduced, we observe that the scattering rates of the individual subbands increase by less than a factor of 2, while the intersubband scattering rate increases by a factor of 7. Such a behavior can be related to the fact that, as the density decreases the energy separation between the two spin-split bands $\Delta_{SO}$ also decreases and therefore it is easier for the carriers to scatter from one subband to another. By reducing the density, the parabolic feature in magnetoresistance around $B=0$T becomes broader and shallower. The decrease of the density from $N=3.0 \times 10^{11}$ to $N=1.7 \times 10^{11}$ causes a reduction of the relative magnetoresistance
Figure 5.7: Left column: Fitting of the low-field magnetoresistivity with the two-band theory [93] (blue lines are data points, and red dashed lines are fitting lines) in the following configurations: (a1) Bochum1282 sample, $V_{TG} = 0V$, $N= 3.0 \times 10^{11} \text{ cm}^{-2}$, (b1) Bochum1282 sample, $V_{TG} = 1V$, $N= 2.3 \times 10^{11} \text{ cm}^{-2}$, (c1) Bochum12029 sample, $N= 3.8 \times 10^{11} \text{ cm}^{-2}$. Right column: Nonlinearity in the low-field Hall resistivity (blue lines are data point, red dashed lines are simulated curves, see text for detail explanations) in the following configurations: (a2) Bochum1282 sample, $V_{TG} = 0V$, $N= 3.0 \times 10^{11} \text{ cm}^{-2}$, (b2) Bochum1282 sample, $V_{TG} = 1V$, $N= 2.3 \times 10^{11} \text{ cm}^{-2}$, (c2) Bochum12029 sample, $N= 3.8 \times 10^{11} \text{ cm}^{-2}$.
(\(\rho_{xx}(0.1T) - \rho_{xx}(0)\))/\(\rho_{xx}(0)\) from 18% to around 7%. This is in agreement with results from (311) p-type heterostructures [94], and represents clear evidence that there is a strong correlation between carrier imbalance in spin-split subbands with different mobilities and positive magnetoresistance. We also observe that an increase of the temperature causes the broadening of the parabolic feature in magnetoresistance around B=0 T, and also an increase of the intersubband scattering rate. By increasing the temperature from 70 mK to 1K at a fixed density of \(3 \times 10^{11}\) cm\(^{-2}\), the relative magnetoresistance \((\rho_{xx}(0.1T) - \rho_{xx}(0))/\rho_{xx}(0)\) drops from 17% to 5%.

It is interesting to point out that the intersubband scattering rate increases faster with increasing temperature than the scattering rates of the individual subbands - this indicates that the presence of the two spin-split subbands in p-type samples might be relevant for the strong temperature dependence of the resistivity of p-type samples even at mK temperatures [94].

Beside the parabolic-like longitudinal magnetoresistance around B=0 T, the presence of the two subbands also modifies the Hall resistivity around B=0T and introduces non-linear corrections [93]. We simulate the Hall resistivity \(\rho_{xy}\) according to eq. (5.1), and use the scattering rates \(K_1\), \(K_2\) and \(K_{12}\), obtained from the fitting of the longitudinal resistivity \(\rho_{xx}\), as input parameters. In order to make this small non-linear corrections to the Hall resistivity visible, we subtract from both, the measured data and the simulated \(\rho_{xy}\), the linear contributions obtained by fitting straight lines to the data and the simulated \(\rho_{xy}\) in the B-field range (-0.15 T, 0.15 T). The non-linear corrections of the Hall resistivity are presented for both, the measured data (blue lines) and simulated \(\rho_{xy}\) (dashed red lines) in Fig. 5.7(a2) (configuration \(V_{TG} = 0\) V, \(N=3.0 \times 10^{11}\) cm\(^{-2}\)) and Fig. 5.7(b2) (configuration \(V_{TG} = 1\) V, \(N=2.3 \times 10^{11}\) cm\(^{-2}\)). We find reasonable agreement between the data and the simulated non-linear corrections of the Hall resistivity.

We further explore the positive magnetoresistance around B=0T in the sample Bochum12029 with a 2DHG 45nm below the surface. A relatively sharp feature around B=0 T is observed, which cannot be fitted with the expression for \(\rho_{xx}\) given in eq. (5.1) if we limit inter-subband scattering \(K_{12}\) only to non-negative values. We found that in the case of \(K_{12}\) being limited to non-negative values, the matching between the data and the fitting curve was closest for \(K_{12} = 0\). This behavior is similar to that observed in anisotropic (311) p-type GaAs samples [94]. If we allow negative values for \(K_{12}\), we can fit the data nicely (Fig. 5.7(c1), blue line is data, dashed red line is fitting curve). The obtained scattering rates of the individual subbands are \(K_1 = 0.10 \times 10^{11}\) s\(^{-1}\), \(K_2 = 0.49 \times 10^{11}\) s\(^{-1}\), while the inter-subband scattering rate is \(K_{12} = -0.072 \times 10^{11}\) s\(^{-1}\). However, the physical meaning of the negative inter-subband scattering rate remains unclear. If we simulate the nonlinear correction of the Hall resistivity in the same way as described above, we obtain larger discrepancies between the data and simulated curve (Fig. 5.7(c2)), than it was the case for the sample Bochum1282.
5.6 Temperature dependence of the SdH oscillations

The temperature evolution of the SdH oscillations in the sample Bochum1282 in the configuration $V_{FG} = 0\,\text{V}$, $N = 3.0 \times 10^{11} \,\text{cm}^{-2}$ is presented in Fig. 5.8. It can be seen in the plot that as the temperature is increased local maxima (minima) in the oscillations related to the lower mobility subband disappear, and at these positions minima (maxima) related to the higher mobility subband remain. This figure shows that the total SdH oscillations are indeed composed of the two sets of oscillations related to the two subbands with different mobilities. It indicates that the oscillations related to the subband with lower mobility (heavier subband mass) weakens faster than those related to the higher mobility subband. However, it should be mentioned that in these measurements a current of 100 nA was applied, and due to stronger sample heating there were not so many oscillations developed as in the measurement with 10 nA (Fig. 5.4).

In order to determine the effective mass of the higher mobility subband (lighter subband mass) we now focus on the oscillations which belong only to this subband and analyze the temperature dependence of the amplitude $\Delta \rho_{xx}/\rho_{xx}$ of the one oscillation around the field $B=0.87\,\text{T}$ (arrow in the left figure) with a Dingle term [26].

![Figure 5.8](image-url)
rate \[26\] and the assumption that carriers from both subbands see the same scattering potential gives \( m_2/m_1 = (\mu_1/\mu_2)^{1/2} \). This relation, together with the previously obtained results for \( m_1, \mu_1 \) and \( \mu_2 \), gives \( m_2 = 0.53 \, m_0 \) for the effective mass of the low-mobility spin-split subband. If we fit a Dingle term to another SdH oscillation around \( B=0.72 \, T \) we obtain a slightly different mass of the higher mobility subband of \( m_1 = (0.30 \pm 0.01) \, m_0 \). This B-field dependence of the effective mass of the spin-split heavy hole subbands is also observed in the experiments by Habib et al. \[95\]. However the values of the effective masses of holes always remain several times larger that the effective mass of electrons in GaAs \( (m_e = 0.067 \, m_0) \).

The obtained values for the effective masses of the two spin-split subbands are comparable to the effective mass of the heavy hole band in bulk GaAs \( (m_h = 0.49 \, m_0) \), which shows that the two spin-split subbands indeed originate from the bulk heavy hole band. However, it should be mentioned here that due to strong non-parabolicity of the heavy hole subband, the effective mass of each of the two subbands depends on the total density in the system \[43, 96\]. Detailed cyclotron resonance studies of the effective masses of the spin-split heavy hole subbands in C-doped (100) GaAs heterostructures can be found in \[68\].

The quantum scattering time in the sample Bochum1282 is obtained by fitting an envelope function to the low-field SdH oscillations, \( \tau_q = 2.3 \times 10^{-12} \, s \). The zero-field longitudinal resistance gives \( \tau_D = 3 \times 10^{-11} \, s \) for the Drude scattering time. The ratio \( \tau_D/\tau_q = 13 \) indicates that a long range scattering potential is dominant in the sample.

Finally it should be mentioned that by increasing the temperature from 70mK to 1K, the zero-field resistivity of the 2DHG in the sample Bochum1282 increases by 6%. This behavior, which has been interpreted as a metallic state of a 2DHG \[85–87\], is another confirmation of the high sample quality.

## 5.7 Evolution of the SdH oscillations in in-plane magnetic field

We also investigate the evolution of the SdH oscillations upon rotation of the sample with respect to the external magnetic field. The dependence of the SdH oscillations on the perpendicular component of the applied B-field is shown in Fig. 5.9. Indicated are the angles between the external B-field and the normal to the sample surface. It can be seen in Fig. 5.9 that the position of the minima and maxima in the SdH oscillations do not change upon increasing the in-plane component of the B-field, which also manifests in no detectable changes in the position of the peaks in the Fourier spectrum of low field SdH oscillations. Therefore, we can conclude that within experimental accuracy, we do not observe changes in the populations of the spin-split subbands upon increasing the in-plane component of the magnetic field. This behavior may indicate that the spin-orbit induced splitting is considerably larger than the Zeeman splitting in (100) 2DHGs. This is not unexpected, since
it is predicted that for GaAs heavy holes with orbital angular momentum $l = 1$ in highly symmetric planes like (100), in-plane g-factor should be 0 [43]. Heavy hole g-factors in (100) p-type GaAs quantum wells have been measured with optically detected magnetic resonance [97], spin-flip Raman scattering [98] and time resolved photoluminescence [99]. The obtained values for the heavy hole g factors in the plane of the quantum well were respectively: less than 0.01 [97], 0 within the experimental accuracy [98], and $0.04 \pm 0.01$ [99]. The measured heavy hole g factors perpendicular to the quantum well were around 2.3 [97, 98].

![Figure 5.9: Dependence of the SdH oscillations on perpendicular component of the applied B-field. Applied current was 100 nA. The angles between the applied field and the normal to the sample surface are indicated. Traces are vertically offset for clarity.](image-url)

**5.8 Summary**

In this chapter we investigated low temperature magnetotransport in two different carbon doped p-type GaAs heterostructures, Bochum1282 and Bochum12029. Observation of the pronounced integer quantum Hall plateaus, features of fractional quantum Hall effect at filling factors 4/3 and 5/3, as well as highly resolved Shubnikov de Haas oscillations documents that C-doped heterostructures have excellent electronic properties. From the measured mobilities of 160'000 cm$^2$/Vs (Bochum1282) and 200'000 cm$^2$/Vs (Bochum12029) at the temperature of 70 mK, we deduce the
mean free path in these 2DHGs to be around 1.5\(\mu\)m (Bochum1282) and 2.0\(\mu\)m (Bochum12029) which is very promising for the fabrication of high quality ballistic p-type nanostructures.

The analysis of the beating of the SdH oscillations reveals the presence of two spin-split subbands which originate from strong spin-orbit interaction in the system. From the measured charge imbalance between these two subbands we deduce that the spin-orbit induced splitting of the heavy hole subband \(\Delta_{SO}\) is around 30\% of the total Fermi energy for both heterostructures. This confirms the presence of extraordinary strong spin-orbit interactions in the investigated C-doped heterostructures.

The observation of the classical positive magnetoresistance around \(B=0\)T in both samples, together with a non-linear correction in the Hall resistivity, also confirms the presence of the two spin-split subbands. Fitting of this Lorentzian-shaped low field magnetoresistance with the two-band transport theory [93] allows the determinations of the scattering rates in individual subbands and their corresponding mobilities, as well as the intersubband scattering rate. From the temperature evolution of the SdH oscillations we extract the effective masses of the spin-split heavy hole subbands \(m_1 = 0.34m_0\) and \(m_2 = 0.53m_0\).
Chapter 6

Weak antilocalization in hole gases

6.1 Introduction - Weak antilocalization in systems with strong spin-orbit coupling

Weak localization is a quantum mechanical effect which arises from the constructive interference between the time reversed partial waves of the charge carriers in disordered materials. This constructive interference leads to an enhanced probability of backscattering of carriers and therefore the longitudinal resistance of the samples increases. Application of a perpendicular magnetic field breaks time reversal symmetry and suppresses weak localization, which manifests as an appearance of a resistance peak in the longitudinal magnetoresistance at B=0 T.

Weak localization arises due to interference of the orbital part of the carriers’ wave-functions. Therefore, the natural question is how the spin degree of freedom of the charge carriers affects weak localization. This question is particularly important in systems with strong spin-orbit (SO) interactions where the spin dynamics is greatly affected by the orbital motion of the charge carriers. It was predicted by Hikami et al. [100] that in the presence of SO interactions the magnetoresistance changes the sign - this effect is called weak antilocalization.

Weak antilocalization is a consequence of the quantum mechanical fact that for the spin 1/2 particle, the spin has to perform a rotation by an angle of $4\pi$ in spin space in order to be mapped into itself. In a diffusive system with SO interaction each scattering event rotates the spin by an infinitesimal angle. These infinitesimal rotations add up along the path to a finite rotation $R$, so that the final spin state is $|s'\rangle = R|s\rangle$, where $|s\rangle$ is the initial spin state of the carrier. For the propagation along the time reversed path the spin experience the same infinitesimal rotations in the opposite direction and opposite sequence and they add up to a finite rotation $R^{-1}$. Therefore the final spin state for the propagation along the time reversed path is $|s''\rangle = R^{-1}|s\rangle$. The spin interference contribution is determined by the matrix element $\langle s'|s''\rangle = \langle s|R^2|s\rangle$. If SO interaction is negligible, the initial spin does not rotate along the paths, and the matrix element $\langle s'|s''\rangle$ is 1, leading to constructive
Chapter 6. Weak antilocalization in hole gases

interference and weak localization. However, if SO interaction is present, the spin experiences the sequence of stochastic scattering events along the path, and upon completing the full path the distribution of the spin directions is isotropic in the absence of external magnetic fields. It was shown by Bergmann [41] that in that case the matrix element \( \langle s'|s'' \rangle \) becomes \(-1/2\) and the spin interference is destructive, leading to weak antilocalization.

Weak antilocalization was first experimentally observed by Bergmann in thin metallic film samples (Fig. 6.1(a)) [101]. The low-temperature magneto-resistance of a thin magnesium (Mg) film covered with a thin layer of gold (Au) has been measured. Mg is a light metal and has very weak SO coupling, while Au has very strong SO coupling. Fig 6.1(a) shows the magneto-resistance of the thin Mg film for different thicknesses of the Au coverage layer (in % of the Mg layer thickness). If there is no Au coverage, which means no SO interaction, the magneto-resistance shows a clear weak localization peak (the lowest trace in Fig. 6.1(a)). As the thickness of the Au layer increases, which corresponds to an increase of the SO interaction strength, a small antilocalization dip superimposed on the larger localization peak starts to develop first, and when the SO interaction becomes strong enough the weak antilocalization dip completely takes over (Fig. 6.1(a)). The ratio between the phase coherence time \( \tau_i \) and SO scattering time \( \tau_{SO} \) is shown on the left side of Fig. 6.1(a) - the stronger the SO interaction, the shorter the time between two SO scattering events \( \tau_{SO} \) is.

Weak antilocalization was subsequently observed also in semiconductor heterostructures [80, 103]. A small antilocalization dip superimposed on a weak localization peak was seen in the magneto-resistance of an inversion layer of InP [103] and n-type GaAs/AlGaAs heterostructure. A fully developed antilocalization dip was observed by Chen et al. in the magneto-resistance of an InAs quantum well [104]. Koga et al. demonstrated the transition from a weak localization peak to a weak antilocalization dip by tunning the symmetry of an InGaAs quantum wells (QW) [102] (Fig. 6.1(b)). They have clearly shown in such a way that the Rashba SO interaction is related to the structural inversion asymmetry of the QW. In the case of a symmetric QW, weak localization is observed, indicating weak SO interactions, while in the case of a highly asymmetric QW, a fully developed antilocalization dip is seen, documenting the presence of very strong SO interactions in that case.

The fitting of the antilocalization data either with Hikami-Larkin-Nagaoka (HLN) theory [100] (SO skew scattering -Elliot mechanism), or Iordanskii-Lyanda Geller-Pikus (ILP) theory [105] (Dyakonov-Perel mechanism), is commonly used for the quantification of the spin-orbit scattering time \( \tau_{SO} \) and spin-orbit strength. It should be mentioned that both theories are developed for diffusive systems, so they are formally applicable only for magnetic fields smaller than the characteristic transport field, \( B < B_tr \), where \( B_tr = h/(2e\ell_m^2) \) and \( \ell_m \) is a mean free path. Besides, these theories consider SO interaction only as a weak perturbation (\( \tau_{SO} \gg \tau_tr \)), and therefore are not fully applicable in the limit of strong SO interactions.

Weak antilocalization is expected to be even more expressed in the case of p-type
6.1. Introduction - Weak antilocalization in systems with strong spin-orbit coupling

Figure 6.1: (a) taken from [101] The magnetoresistance of a thin Mg film at 4.5K for different coverages with Au (shown in % of the Mg layer thickness on the right side). As the SO coupling increases a weak antilocalization dip develops. (b) taken from [102] Low field magnetoresistance of InGaAs quantum wells (QW) with different well symmetry. In the case of symmetric QW weak localization peak is present (lowest trace), and as the symmetry of the QW is changed to more asymmetric, weak antilocalization dip develops (the topmost trace). The background resistance in these samples was in the range 1200-1600Ω.

GaAs heterostructures due to strong SO interactions in these systems. Experimental studies of weak antilocalization in Be-doped p-type GaAs heterostructures are reported in [106, 107].

In this chapter we report measurements of the weak antilocalization in two C-doped p-type GaAs heterostructures (Bochum1282 and Bochum12029). The observation of a fully developed antilocalization dip in the magnetoresistance of both heterostructures clearly demonstrates the presence of very strong spin-orbit interactions. By fitting the antilocalization dip by HLN theory in a narrow B-field range around B=0 T we extract the phase coherence time of holes at \( T = 70 \) mK to be 190 ps in the sample Bochum1282 and 340 ps in sample Bochum12029. We investigate the temperature dependence of the phase coherence time and find that it reasonably satisfies a \( 1/T \) dependence. In the sample Bochum1282, which is equipped with a top-gate, we investigated the density dependence of antilocalization. The antilocalization dip was present for all accessible densities in the range \( 2 - 3 \times 10^{11} \) cm\(^{-2}\), and we observed a decrease of the phase-coherence time as the density is reduced. Due to the high mobility of the investigated samples, the transport field \( B_{tr} \) is less than 0.3 mT in both samples, and fitting with the HLN theory is formaly valid only for B-fields smaller than this value. Due to the strong SO interactions, the characteristic SO field \( B_{SO} = \hbar/(4D\tau_{SO}) \) is much larger than \( B_{tr} \), and therefore the fitting of the high-field tails of the antilocalization dip and the extraction of \( \tau_{SO} \) from the
data is limited. We also discuss the limitations of an accurate determination of the tails of the antilocalization dip, due to the presence of other contributions in the low-field magnetoresistance of 2DHGs, like the classical Lorentzian background resistance (see chapter 5) or the background resistance due to carrier-carrier Coulomb interactions [107–109].

6.2 Weak antilocalization correction

Weak (anti)localization effects are, in principle, better visible in lower mobility samples in which the carrier mean free path is much smaller compared to the phase-coherence length. In higher mobility samples where \( k_F l_m \gg 1 \) localization effects are weaker and harder to resolve compared to low mobility diffusive systems. For the two high mobility p-type GaAs heterostructures investigated in this chapter \( k_F l_m \sim 100 \sim 300 \), and therefore the magnitude of the localization effects was expected to be very small. In order to resolve an anti-localization peak in the magnetoresistivity of these two samples we had to measure both, the resistance and the magnetic field with very high accuracy.

The two investigated heterostructures (Bochum1282 and Bochum12029) are patterned into Hall-bars with a width of \( W = 100\mu m \) and a separation between the contacts \( L = 500\mu m \). We performed four-point measurements of the resistivity using standard lock-in techniques. A current of 20 nA is applied to the sample at a frequency of 31 Hz, and the voltage is measured with an integration time of 300 ms. In order to reduce the noise we used a voltage amplifier with an amplification of 1000 directly at the outputs from the cryostat. Besides, to further increase the signal to noise ratio around the antilocalization peak we averaged over 25 measurements at each magnetic field, with an interval of 1.5 s between measurements. In this way we managed to reach a noise level of less than 0.03 Ω for measured resistances above 200 Ω. A special, home-built power supply for the magnet is used and the magnetic field is measured with a resolution of 40 \( \mu T \) in the case of the measurements on the sample Bochum1282 and 20 \( \mu T \) in the case of the measurements on the sample Bochum12029.

In the left column of Fig. 6.2, the raw magnetoresistivity data are presented for the top-gated sample Bochum1282 in the gate-configurations \( V_{TG}=1V, N = 2.3 \times 10^{11} \text{ cm}^{-2}, \mu = 130'000 \text{ cm}^2/\text{Vs} \) (Fig. 6.2 (a1)) and \( V_{TG}=0V, N = 3 \times 10^{11} \text{ cm}^{-2}, \mu = 160'000 \text{ cm}^2/\text{Vs} \) (Fig. 6.2 (b1)), as well as for sample Bochum12029 with \( N = 3.8 \times 10^{11} \text{ cm}^{-2}, \mu = 200'000 \text{ cm}^2/\text{Vs} \) (Fig. 6.2 (c1)). In all three cases we observe a sharp antilocalization resistance dip around B=0T. For the higher mobility sample Bochum12029 with \( k_F l_m = 315 \), the weak antilocalization dip has a magnitude \( \sim 0.05\Omega \) and a width of \( \sim 0.3\text{mT} \) (Fig. 6.2 (c1)). In the sample Bochum1282 (Fig. 6.2 (a1),(b1)) we observe that as the factor \( k_F l_m \) decreases due to a reduction of the sample mobility and density with a top gate, the magnitude and the width of the antilocalization dip become larger.
6.2. Weak antilocalization correction

Figure 6.2: **Left column** Raw magnetoresistivity data at $T=65$ mK: (a1) sample Bochum1282, $V_{TG}=1V$, $N = 2.3 \times 10^{11}$ cm$^{-2}$, $\mu = 130'000$ cm$^2$/Vs; (b1) sample Bochum1282, $V_{TG}=0V$, $N = 3\times 10^{11}$ cm$^{-2}$, $\mu = 160'000$ cm$^2$/Vs, (c1) sample Bochum12029, $N = 3.8 \times 10^{11}$ cm$^{-2}$, $\mu = 200'000$ cm$^2$/Vs; **Right column** Quantum correction of the resistivity obtained after subtraction of the classical 2-band positive magnetoresistivity (see chapter 5) from the raw data: (a2) sample Bochum1282, $V_{TG}=1V$, (b2) sample Bochum1282, $V_{TG}=0V$, (c2) sample Bochum12029.
Chapter 6. Weak antilocalization in hole gases

It was discussed in chapter 5 that due to the presence of the two spin-split heavy hole subbands, a classical Lorentzian magnetoresistance develops around $B=0$. In order to separate the quantum correction from the low-field magnetoresistance, we subtract the classical positive magnetoresistivity $\rho_{\text{class}}$ from the total resistivity $\rho$. These differences are plotted in the right column of Fig. 6.2. The classical contribution is obtained by fitting the data in the range $(-0.25T, -0.01T) \cup (0.01T, 0.25T)$ with the two-band Zaremba model discussed in chapter 5. We chose the upper limit ($\pm 0.25T$) where SdH oscillations start to develop and a lower limit ($\pm 0.01T$) where quantum corrections are expected to be suppressed.

As it can be seen in Fig. 6.2(a2,b2,c2), in all three cases only a weak antilocalization dip is present in the low-field magnetoresistance. The fact that a fully developed antilocalization dip is observed, and no superposition of an antilocalization dip and a larger localization peak, confirms that the spin-orbit interactions in the system are exceptionally strong [101, 110].

### 6.3 Phase-coherence time of holes

In order to proceed with fitting of the data with HLN-theory, we need to calculate the conductivity correction

$$\Delta \sigma(B) = (\sigma(B) - \sigma(0)) - (\sigma_{\text{class}}(B) - \sigma_{\text{class}}(0)), \quad (6.1)$$

where $\sigma$ is the longitudinal conductivity, obtained from the inversion of the measured resistivity tensor, and $\sigma_{\text{class}}$ is the classical longitudinal conductivity, obtained from the fitted $\rho_{\text{class}}$. The obtained conductivity correction $\Delta \sigma(B)$ is plotted in Fig 6.3(a) for sample Bochum1282 and Fig. 6.3(b) for sample Bochum12029. In Fig. 6.3 the dots represent the measured data and the full lines are fits with the HLN-theory. As it was mentioned before HLN-theory is valid in the diffusive regime, where $B < B_{tr} = \hbar/(2e\ell_m^2)$. In the case of sample Bochum1282 $B_{tr} < 0.3$ mT, while for the higher mobility sample Bochum12029 $B_{tr} < 0.1$ mT. We took the fitting regions slightly larger than these values (shown in Fig. 6.3) in order to have a reasonable number of points to perform the fitting. The data in the given B-field ranges are fitted with the expression [100, 108]

$$\Delta \sigma(B) = -\frac{e^2}{\pi \hbar} \left( \frac{1}{2} \Psi \left( \frac{1}{2} + \frac{B_{\varphi}}{B} \right) - \frac{1}{2} \ln \frac{B_{\varphi}}{B} \right), \quad (6.2)$$

where $\Psi(x)$ is the digamma function, $B_{\varphi} = \hbar/(4De\tau_{\varphi})$, $D$ is the diffusion constant and $\tau_{\varphi}$ is the phase-coherence time. The only fitting parameter is $B_{\varphi}$. Satisfactory fitting is obtained (full lines in Fig. 6.3) for both samples and the phase coherence time of holes is extracted. For the sample Bochum1282 (Fig. 6.3(a)) in the configuration $V_{TG} = 1$V, $k_Fl_m = 120$ (red points) $B_{\varphi} = 5.1 \times 10^{-5}$T, $\tau_{\varphi} = 165$ps, and in the configuration $V_{TG} = 0$V, $k_Fl_m = 200$ (black points) $B_{\varphi} = 2.6 \times 10^{-5}$T,
6.4. Temperature dependence of the weak antilocalization correction

\[ \tau_\varphi = 190\text{ps} \]. For the sample Bochum12029 (Fig. 6.3(b)) with \( k_F l_m = 315 \) the obtained values are \( B_\varphi = 0.95 \times 10^{-5}\text{T} \) and \( \tau_\varphi = 340\text{ps} \). The corresponding phase coherence lengths, calculated according to the expression \( l_\varphi = \sqrt{D/\tau_\varphi} \), valid for the diffusive regime, are 1.8 \( \mu \text{m} \), 2.5 \( \mu \text{m} \) and 4.1 \( \mu \text{m} \), respectively. These values are compatible with the values obtained from measurements of Aharonov-Bohm oscillations (see chapter 9), and demonstrate that the fabrication of phase-coherent p-type GaAs nanostructures is accessible with present nanofabrication technologies. The obtained values for the holes’ phase coherence time show that it decreases as the density in the sample is reduced. Since Coulomb interactions become stronger at lower densities, this tendency suggests stronger charge dephasing of holes when the carrier-carrier interactions become more prominent.

![Figure 6.3: Fitting of the antilocalization conductance peak with HLN-theory eq. (6.2)-full lines are fitting curves, and points are experimental data for: (a) Sample Bochum1282 in configurations \( V_{TG} = 1\text{V} \), \( k_F l_m = 120 \) (red line) and \( V_{TG} = 0\text{V} \), \( k_F l_m = 200 \) (black line). (b) Sample Bochum12029, \( k_F l_m = 325 \) (blue line)](image)

### 6.4 Temperature dependence of the weak antilocalization correction

In the remaining part of this chapter we show only the data from the sample Bochum1282 in the configuration \( V_{TG} = 1\text{V} \), \( k_F l_m = 120 \), since in that configuration the antilocalization conductance peak is most strongly visible. Fig. 6.4 shows the temperature evolution of the resistivity around \( B=0\text{T} \). The antilocalization dip depends strongly on temperature and disappears completely above 300 mK. It can also be seen that the sample exhibits metallic behavior in its temperature dependence, with the zero-field resistivity increasing with temperature, which confirms the high sample quality.

The fitting of the conductance antilocalization peak is performed for each measured temperature and the phase-coherence times are extracted. From Fig. 6.5 it
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Figure 6.4: (a) Temperature dependence of the antilocalization resistivity dip in the sample Bochum1282 with top-gate $V_{TG} = 1\text{V}$

can be seen that the inverse phase-coherence time $1/\tau_\phi$ depends almost linearly on temperature.

Figure 6.5: (a) Fitting of the antilocalization conductance peak with HLN-theory eq. (6.2) for temperatures 70 mK and 190 mK - full lines are fitting curves, and points are experimental data. (b) Temperature dependence of the inverse phase-coherence time.

6.5 Limitations in extracting spin-orbit scattering times

Before we proceed with the evaluation of the spin-orbit scattering time $\tau_{SO}$ from antilocalization measurements we will estimate $\tau_{SO}$ from Shubnikov-de Haas mea-
6.5. Limitations in extracting spin-orbit scattering times

The estimated spin-orbit induced splitting of the heavy hole band at a density of \( N = 2.3 \times 10^{11} \text{ cm}^{-2} \) is \( \Delta_{SO} = 0.47 \text{meV} \) (Fig. 5.6). In semiconductor heterostructures with inversion asymmetry the dominant spin-orbit relaxation mechanism is the Dyakonov-Perel mechanism [111], which establishes the following relation between the spin-orbit (SO) induced splitting of the band \( \Delta_{SO} \) and SO relaxation time \( \tau_{SO} \):

\[
\tau_{SO}^{-1} = \Delta_{SO}^2 \frac{\tau_{tr}}{4\hbar^2} \quad [111].
\]

Knowing \( \tau_{tr} = 26 \text{ps} \) we estimate \( \tau_{SO} \sim 0.3 \text{ps} \). It is clear that the SO interactions are so strong, so that \( \tau_{SO} \ll \tau_{tr} \), and therefore the SO interaction cannot be treated only as a weak perturbation \( \tau_{SO} \gg \tau_{tr} \), which is the assumption of all present theories. If we further estimate the characteristic field \( B_{SO} = \frac{\hbar}{4De\tau_{SO}} \), at which the effects of SO interactions become suppressed and the weak antilocalization positive magnetoresistance turns into a weak localization negative magnetoresistance [112] we obtain \( B_{SO} \sim 30 \text{mT} \), which is far beyond the transport field \( B_{tr} \sim 0.3 \text{mT} \) up to which diffusive theories of weak antilocalization are applicable. However, the value of \( B_{SO} \sim 30 \text{mT} \) provides a qualitative understanding of the fact that we observe only a weak antilocalization dip without a weak localization peak in our data.

We now perform the fitting of the data with the HLN-theory with SO-interactions included using the expression [100, 105, 112]:

\[
\Delta\sigma(B) = -\frac{e^2}{\pi\hbar} \left( \frac{1}{2} \Psi\left(\frac{1}{2} + \frac{B_{\varphi} + B_{SO}}{B}\right) - \frac{1}{2} \Psi\left(\frac{1}{2} + \frac{B_{\varphi} + 2B_{SO}}{B}\right) 
- \frac{1}{2} \ln \frac{B_{\varphi}}{B} + \ln \frac{B_{\varphi} + B_{SO}}{B} + \frac{1}{2} \ln \frac{B_{\varphi} + 2B_{SO}}{B} \right) \quad (6.3)
\]

It should be mentioned that the HLN-theory was originally developed for metallic samples where the Elliot SO skew-scattering mechanism is present. For this mechanism the spin-splitting energy is proportional to \( k^3 \), where \( k \) is the wavevector. However, in most semiconductor heterostructures the Dyakonov-Perel spin relaxation is dominant [111]. The theory by Iordanski-Lyanda Geller-Pikus (ILP) describes the weak antilocalization correction for this type of spin relaxation, and it involves both linear and cubic in \( k \) spin-splitting terms [105]. However, it is shown that if the linear contribution is negligible, and only cubic spin-splitting is present, ILP-theory gives the same result as HLN-theory [105, 112], expressed with equation eq. (6.3). Since the spin-orbit induced splitting of the heavy hole band is proportional to \( k^3 \) [43, 49], it is suitable to use the HLN equation eq. (6.3) for fitting the antilocalization in hole systems. The given equation contains two fitting parameters, namely \( B_{\varphi} \) and \( B_{SO} \).

The sharpness of the antilocalization conductance peak is mostly determined by \( \tau_{\varphi} \), while the tail of the peak depends on \( \tau_{SO} \) [108, 112]. Therefore we investigate in Fig. 6.6(a) the fitting of the data with equation eq. (6.3) in two different magnetic field ranges: (-0.5mT, 0.5mT) (blue line) and (-5mT, 5mT) (red line). It can be seen that fitting in the narrow range (-0.5mT, 0.5mT) reproduces quite well the low field behavior up to \( B \sim B_{tr} \), but above that field the fit does not match the
experimental data. The obtained fit parameters are $\tau_\varphi = 210\text{ps}$ and $\tau_{SO} = 40\text{ps}$. On the other hand the fit of the data in the larger range (-5mT, 5mT) matches better the tails of the peak, but does not fit satisfactorily the low-field data below $\sim B_{tr}$. The fit parameters in this second case are $\tau_\varphi = 120\text{ps}$ and $\tau_{SO} = 3\text{ps}$. While the obtained values for $\tau_\varphi$ differ by less than a factor of 2, and are comparable with the value obtained by fitting with the equation eq. (6.2) which neglects the contribution from $\tau_{SO}$, the obtained values for $\tau_{SO}$ differ by more than an order of magnitude. While the fitting in the range $(-0.5\text{mT}, 0.5\text{mT})$, below $B_{tr}$ is more justified from the formal point of view, it is clear that it underestimates the SO strength, since it gives an upturn from antilocalization to localization which is not present in the data. Also if we compare the $\tau_{SO} = 40\text{ps}$, obtained from the fit in the range $(-0.5\text{mT}, 0.5\text{mT})$, with that estimated from the beating of SdH oscillations $\tau_{SO} = 0.3\text{ps}$ we see a huge discrepancy. Fitting in the range $(-5\text{mT}, 5\text{mT})$ gives closer agreement between the extracted $\tau_{SO} = 3\text{ps}$, and that obtained from SdH oscillations. In Fig 6.6(b) we investigate influence of changing $\tau_{SO}$ at fixed $\tau_\varphi = 210\text{ps}$ on the fitting curves, and find that the fitting procedure becomes less sensitive for $\tau_{SO}$ smaller than 3ps. Therefore, rather than giving the exact value, this fit sets the upper limit on $\tau_{SO}$.

We have also tried to fit the data with the theory by Averkiev et al. [113], which considers anisotropic spin relaxation and involves 3 instead of 2 free parameters, but could not obtain a better fit in the whole data range.

To summarize, we have observed that the measured weak antilocalization correction cannot be satisfactorily fitted in the whole measurement range, and therefore the estimation of the spin relaxation time from this measurement is limited. From the theoretical point of view there might be two reasons for such an observation:

- Spin-orbit interactions are very strong in the investigated systems and cannot be
6.5. Limitations in extracting spin-orbit scattering times

considered only as a perturbation, as it is taken in all present theories. Similar difficulties in fitting the weak antilocalization data are observed for InGaAs/InP quantum well with extremely strong SO interactions [114].

- The mobility in the systems is very high, and this requires a weak antilocalization theory for non-diffusive systems. Such theories exist [115–117], but contain complex expressions which are difficult to fit to the data.

![Figure 6.7: Quantum conductivity corrections obtained after subtraction of the classical background in different B-field ranges indicated in the figure.](image)

The observed difficulties in fitting the data with theoretical expressions for weak antilocalization might also originate from the presence of some other contributions in the low-field magnetoresistance. First, the determination of the classical background, which is subtracted from the data is not unique. Although the choice of the range where the fitting of the classical background is performed is supported by physical reasoning, the borders of the fitting range could have also been chosen slightly different. In Fig. 6.7 we analyze the differences in the quantum corrections of the conductivity depending on the choice of the B-field range where the fitting of the classical background was performed (indicated in Fig. 6.7). We see that the conductance peak itself is not affected by the choice of fitting background up to ∼ 15 mT, and only beyond this field the tails of antilocalization peaks become dependent on the subtracted classical background. Therefore, this is probably not the reason for the observed discrepancies in the fitting procedures.

It is also possible that the low-field magnetoresistance contains some contributions due to carrier-carrier Coulomb interactions [63, 108, 109]. These interaction-corrections might be particularly strong in p-type GaAs systems due to the larger effective mass of holes.
6.6 Summary

We have observed weak antilocalization in two investigated C-doped GaAs heterostructures. The observation of only a weak antilocalisation resistance dip, instead of a transition from an antilocalization dip to a localization peak confirms the presence of strong spin-orbit interactions in the system. By fitting the weak antilocalization data in the B-field range up to $B \sim B_{tr}$ we extract the phase coherence times of holes to be around 190ps for sample Bochum1282, and 340ps for sample Bochum12029. The temperature dependence reveals that the antilocalization dip persists up to around 300mK and that $1/\tau_\varphi$ depends on temperature in an almost linear fashion. This indicates that carrier-carrier Coulomb interactions are the dominant dephasing mechanism. The extraction of spin orbit relaxation time from fitting of the antilocalization data appeared to be limited, due to the fact that the sample mobility is high and the spin-orbit interactions are strong and cannot be treated as a perturbation.
Chapter 7

Ballistic transport in p-type Quantum Point Contacts

7.1 Introduction

Conduction quantization, discovered in 1988 by van Wees et al. [31] and Wharam et al. [32] is the main characteristic of ballistic transport in clean one-dimensional systems. Such systems can be considered as electron (hole) waveguides, where each spin-degenerate mode carries one conductance quantum $2e^2/h$. The step-wise conductance changes in one dimensional systems can be well explained within the single-particle picture: the total conductance $N \times 2e^2/h$ is determined by the integer number of one dimensional (1D) subbands $N$ below the Fermi level.

Although a material independent phenomenon, conductance quantization was so far mostly studied in quantum point contacts (QPC) fabricated on n-type GaAs heterostructures, due to their high electronic quality. However, it is known that in n-type GaAs systems neither electron-electron nor spin-orbit interactions are strong. Therefore, the investigation of the conductance quantization in the limit of strong carrier-carrier Coulomb interactions, where the single particle picture is not valid anymore, or in the limit of strong spin-orbit interaction where different modes can be mixed, may give new insights in this effect.

Beside plateaus at integer multiples of $2e^2/h$, in many experiments an additional plateau-like feature around $0.7 \times 2e^2/h$ has been observed [118]. The understanding of this anomaly has been puzzling physicists during the last decade. There have been several theoretical proposals, all of them invoking many-body effects, but none of them could fully explain observed experimental results. Recently, Rejec and Meir presented numerical spin density-functional calculations [119] that reveal the formation of a quasibound electronic state with a spin-1/2 magnetic moment in the QPC as its conductance rises towards the first plateau. This local quasibound spin state formed within a QPC can be screened by the spins from the Fermi sea in the leads and therefore Kondo-like effects may be important for understanding the 0.7 feature. The observed temperature and magnetic field dependence of the 0.7 feature
in the experiments by Cronenwett et al. [120] indeed point to the Kondo origin of the 0.7 feature. It is interesting to mention that the theory of Rejec and Meir also predicts, although less robust, an anomalous shoulder around $1.7 \times 2e^2/h$ which is so far observed only in very few experiments [121].

Carrier-carrier Coulomb interactions, as well as spin-orbit interactions are stronger in p-type GaAs heterostructures compared to their n-type counterparts and therefore p-type QPCs provide fertile ground for the investigation of these interaction effects in 1D transport. However, due to technological difficulties to fabricate stable p-type GaAs nanodevices with conventional split-gate techniques, transport phenomena in p-type QPCs remained largely unexplored. Ballistic transport of holes was first demonstrated in magnetic focusing experiments by Heremans et al. in 1992 in a device fabricated by wet-etching from Si-doped (311) heterostructure [59]. Conductance quantizations in a p-type QPC was observed for the first time by Zailer et al. in 1994 in a device fabricated by wet etching [21]. This structure suffered from time dependent fluctuations in the measured signal, and it was necessary to perform averaging of 40 conductance traces in order to resolve the conductance plateaus. Later, conduction quantization was observed in single conductance traces in p-type QPCs defined by metallic split-gates [71] and AFM oxidation lithography [23], although still with limited quality. Significant improvement in the quality of the conductance quantization was recently obtained in specially designed p-type 1D structures. Danneau et al. measured conductance quantization in a hole quantum wire formed by surface gates on top of a 2D bilayer structure, and for the first time observed the anomalous 0.7 plateau in a hole system [122]. However, it remains unexplained why the stability of these devices fabricated in p-type bilayer structure is significantly improved compared to conventional devices defined on single-layer structures. Klochan et al. reported very stable conductance plateaus together with a 0.7 feature in an induced one-dimensional hole system fabricated on an undoped GaAs/AlGaAs heterostructure [123]. They argue that the absence of remote doping in their structure is responsible for the improved quality of the conductance quantization. Finally, Rokhinson et al. investigated the 0.7 structure in an AFM defined p-type QPC by employing spatial spin separation in magnetic focusing experiments [124], and found that holes injected from a QPC set around $0.7 \times 2e^2/h$ are spin polarized [125]. This static spin polarization of holes injected from a QPC questions the Kondo interpretation of the 0.7 anomaly, which requires dynamic spin polarization in a QPC.

In this chapter we present conductance measurements in a QPC fabricated by AFM oxidation lithography on a modulation doped p-type (100) GaAs heterostructure with a hole gas 45 nm below the surface (wafer Bochum12029). The observation of the plateaus at $1 \times 2e^2/h$ and $2 \times 2e^2/h$ demonstrates ballistic transport of holes through the QPC. Besides, we observe a plateau-like structure at $0.8 \times 2e^2/h$ at 500 mK, which evolves into a dip-like structure below the first plateau as the temperature is reduced to 70 mK. We also find plateau like feature at $1.7 \times 2e^2/h$ at the temperature of 70 mK. In differential conductance vs. bias measurements we
observe a pronounced zero-bias peak for a QPC conductance $\sim 0.8 \times 2e^2/h$, which weakens as the conductance increases to $2e^2/h$, and completely disappears above the first plateau. This behavior might indicate that the structure below the first plateau is related to Kondo-like effect [120]. All features observed in this sample were stable and reproducible in several different cool-downs.

7.2 Conductance quantization and 0.7 feature in a p-type QPC

The sample was fabricated by AFM oxidation lithography on a p-type carbon doped (100) GaAs heterostructure, with a shallow 2DHG located 45 nm below the surface. An AFM micrograph of the QPC is shown in inset of Fig. 7.1. The lithographic width of the QPC is $\sim 165$ nm. The hole density in an unpatterned sample is $3.8 \times 10^{11}$ cm$^{-2}$ and the mobility is 200 000 cm$^2$/Vs at a temperature of 70 mK, corresponding to a Fermi wavelength of $\sim 40$ nm, and a mean free path of $\sim 2 \mu$m.

We have measured the differential conductance of the QPC by applying an ac source-drain bias of 10 $\mu$V symmetrically across the QPC and measuring the current in a two-terminal configuration with lock-in techniques. The measurements were performed in a dilution refrigerator with temperatures in the range from 70 mK to 500 mK.

In Fig. 7.1 we present the conductance of the QPC at temperatures of 70 mK (upper trace) and 500 mK (lower trace). The fact that the QPC closes as the plunger gate voltage increases to more positive values clearly demonstrates that we measure hole transport. The two presented curves are obtained after subtracting the background resistance of 12.2 k$\Omega$ from the raw data. The value of the background resistance is determined in such a way that the first conductance plateau is well aligned with $1 \times 2e^2/h$. The background resistance consists of two parts: the contact resistance and the resistance of the 2DHG regions around the QPC. The measured two-terminal contact resistances in the unpatterned mesa structure were around 10 k$\Omega$, consistent with the value of the subtracted background resistance.

At the temperature of 500 mK we observe clear conductance plateaus at $1 \times 2e^2/h$ and $2 \times 2e^2/h$ (black arrows in Fig 7.1) together with a plateau-like feature at $0.8 \times 2e^2/h$ (gray arrow in Fig. 7.1). As the temperature is lowered to 70 mK the plateau-like structure at $0.8 \times 2e^2/h$ evolves into a dip-like structure with its minimum remaining around $0.8 \times 2e^2/h$ and its maximum approaching $1 \times 2e^2/h$. Such a behavior agrees with earlier observations that the 0.7 plateau is stronger at higher temperatures, and that upon reducing the temperature it evolves towards $2e^2/h$ plateau [118, 120]. We also find that upon reducing the temperature to 70 mK another plateau-like structure develops around $1.7 \times 2e^2/h$, consistent with the prediction of Rejec and Meir [119]. Beside these broad conductance features we can see several sharp resonances on the traces in Fig. 7.1. Within one cool-down these sharp resonances remain at fixed plunger gate values for all plunger gate sweeps,
and they remain visible even when the temperature is increased to 500 mK (lower trace in Fig. 7.1). The sample shows very good stability in time and we do not observe any hysteretic behavior in plunger gate sweeps.

In order to check the reproducibility of the observed conductance features we performed another cool-down of the same sample in a different dilution fridge, several months after the first measurements. The conductance curves from the two different cool-downs are shown in Fig. 7.2. While the positions and shapes of the sharp resonances change from one cool-down to another, the broad conductance features are reproducible in different cool-downs and remain at the same plunger gate voltages. This indicates that the sharp resonances are related to the particular configuration of the frozen impurities, which changes from one cool-down to another, while the broad conductance features really originate from the QPC confining potential.

Beside the discussed Kondo-related origin, we cannot exclude other possible mechanisms being responsible for the observed conductance features. Similar dips in the conductance of quantum wires are predicted to occur in the presence of strong spin-orbit interaction due to SO induced mixing of different modes [126]. In Ref. [127] it was discussed that in the presence of boundary roughness scattering the conductance of quantum wires might show broad dips between the adjacent plateaus. However it is known that AFM defined oxide lines create relatively smooth potential
7.3. Nonlinear differential conductance measurements

Source-drain bias spectroscopy is commonly used for the determination of the energy separation between the subsequent transverse modes in a QPC [129, 130]. In Fig. 7.3(a) we present the measurements of the QPC differential conductance $G_{\text{diff}} = dI/dV$ as a function of applied dc source drain bias and a QPC plunger gate voltage $V_{\text{PG}}$. In addition to a dc source-drain bias, a small ac source-drain bias of 10 $\mu$V is applied symmetrically across the QPC and the ac current is measured. By changing the plunger gate voltage in the range from 355 mV to -345 mV we are able to tune the QPC from pinch-off to a conductance of around $3 \times 2e^2/h$.

We observe in Fig. 7.3(a) an accumulation of traces around $1 \times 2e^2/h$ and $2 \times 2e^2/h$, with additional structure developing just below these plateaus around zero bias. However, at larger biases we do not observe the expected accumulation of conductance traces at half-plateaus, possibly due to the presence of stronger noise...
Figure 7.3: (a) Differential conductance of the QPC as a function of source-drain bias for plunger gate voltages $V_{pg}$ in the range from 355 mV (bottom-most curve) to -345 mV (top-most curve). (b) QPC transconductance as a function of source-drain bias and plunger gate voltage $V_{pg}$. White regions correspond to low and black regions to high transconductance. Transconductance is obtained by numerical differentiation of the differential conductance with respect to the $V_{pg}$. (c) Differential conductance as a function of applied bias for plunger gate voltages 60 mV, 110 mV, 210 mV and 310 mV. Zero bias peak is present below and absent above the first plateau.

In order to estimate the sub-band level spacing we plot in Fig. 7.3(b) the QPC transconductance $dG_{diff}/dV_{PG}$ as a function of the applied dc bias. The transconductance is obtained by taking the derivative of the differential conductance data with respect to $V_{PG}$. Before taking the derivative we performed smoothing of the data, to suppress a further increase of the noise by the derivative. White regions correspond to low transconductance (plateaus in conductance), while black regions correspond to high transconductance. We find well defined diamond-like structure in the transconductance for the second plateau, while the structure around the first plateau is not so clearly developed. From the extent of the diamond in bias direction of $\sim 0.8$ mV, we can estimate the sub-band spacing between the 2nd and 3rd QPC mode. However, we have to take into account that the bias of 0.8 mV is a total bias, spread between the QPC and Ohmic contacts. Knowing that the contact resistance is $R_{cont} = 12.2$ k\ohm, and taking for the resistance of the QPC at the second plateau $R_{QPC} = 6.5$ k\ohm, we estimate the voltage drop across the QPC to be...
7.4 Evolution of plateaus in in-plane magnetic field

\[ V_{QPC} = V_{bias} \times \left( \frac{R_{QPC}}{R_{QPC} + R_{cont}} \right) \approx 0.3 \text{ meV}. \]

If we use a simple expression for the separation of the energy levels in a 1D infinite potential well \( \Delta E = \frac{\pi^2 \hbar^2}{2m^*l^2} \), from the estimated sub-band separation of 0.3 meV we estimate the electronic width of a QPC to be around 60 nm. This agrees with the fact that the lithographic width of a QPC is 165 nm and the depletion length of oxide lines is typically around 100 nm [75].

Fig. 7.3(c) shows how differential conductance vs. bias traces evolve as the QPC conductance is tuned by the plunger gate from values below, to values above \( 2e^2/h \). We observe that below the first plateau (\( V_{PG} = 310 \) mV) the strong zero-bias peak develops in differential conductance. As the QPC conductance approaches the first plateau the zero-bias peak gets weaker (\( V_{PG} = 210 \) mV, \( V_{PG} = 110 \) mV) and above the first plateau (\( V_{PG} = 60 \) mV) completely disappears. The presence of a zero-bias peak in differential conductance indicates a possible signature of Kondo effect. We observe that the zero-bias peak weakens as the temperature is increased, but is still present at 500 mK. We also find that the zero bias peak disappears upon applying an in-plane magnetic field.

7.4 Evolution of plateaus in in-plane magnetic field

We explore in Fig. 7.4 the evolution of the QPC conductance in in-plane magnetic field. Fig. 7.4(a) shows the conductance, with constant background of 12.2 kΩ subtracted, for in-plane magnetic field stepped from 0 to 12 T in steps of 1T. However, it is known that strong in-plane magnetic fields can cause the depopulation of spin-split subbands, and change the background resistance of the sample. The magnetic field dependent background resistance of the sample, shown in the inset of Fig. 7.4(a), is determined in such a way that the lowest plateau in the conductance is always aligned with \( 2e^2/h \) after subtracting this background. The data obtained after subtracting this magnetic field dependent background are presented in Fig. 7.4(b).

We observe that the \( 2 \times 2e^2/h \) and \( 1 \times 2e^2/h \) plateaus become weaker and eventually disappear as the in-plane B-field increases. The plateau-like structure around \( 2.5 \times 2e^2/h \) starts to develop at B-fields larger than 8T. However, we do not see a clear transition, where both even and odd plateaus are present, like in the case of n-type point contacts. We also find that the dip-like structure below the first plateau evolves downwards, towards \( 0.5 \times 2e^2/h \) as the in-plane B-field is increased, in agreement with the expected behavior of the 0.7 structure [118].

7.5 Summary

We have measured the conductance of an AFM defined QPC on a p-type GaAs heterostructure. The observation of clear conductance plateaus at \( 1 \times 2e^2/h \) and \( 2 \times 2e^2/h \) provides evidence for ballistic transport of holes in the structure. We also
find a signature of the so-called 0.7 anomaly, by observing plateau-like structure in the conductance around $0.8 \times 2e^2/h$ at a temperature of 500 mK. This plateau-like structure evolves into a dip-like structure with its minimum remaining around $0.8 \times 2e^2/h$ and its maximum approaching $1 \times 2e^2/h$ as the temperature is lowered to 70 mK. The observed zero-bias peak in differential conductance of a QPC set below the first conductance plateau suggests that the structure below the first plateau originates from a Kondo-like effect. However, based on measurements on this one sample, we are not able to eliminate other possible mechanisms being responsible for the observed conductance features. The investigation of QPCs with different geometries and more adiabatic confining potentials is necessary in order to gain
clearer understanding of the 0.7 structure.
Chapter 8

Single hole transistors in p-type GaAs heterostructures

8.1 Introduction

Semiconductor quantum dots are small conducting islands fabricated in semiconductor structures, that can have a tunable number of confined charge carriers. The electronic states of the confined carriers, as well as their coupling to the leads can be well controlled with external gate voltages. Because quantum dots can contain individual charges and spins, they are envisioned to be very promising candidates for the experimental realization of quantum computation and therefore they attracted considerable research interest recently. However, it is not the charge, but rather the spin degree of freedom of the confined carriers that is considered as a potential carrier of the fundamental unit of information [131]. This is because spins are less affected by the environmental noise than charges, due to their weaker coupling to the surrounding. In order to be able to perform quantum computation with quantum dot spins, one needs to manipulate them on a time scale faster than spin relaxation time $T_1$ and spin decoherence time $T_2$. Two dominant spin relaxation and spin decoherence mechanisms in quantum dots are spin-orbit interaction [132] and hyperfine interaction with nuclear spins [133]. However it has been shown that for all spin-orbit mechanisms the spin decoherence times in quantum dots are as large as the spin relaxation times ($T_2 = 2T_1$) [134] at low temperatures ($k_B T \ll \Delta$, $\Delta$ being the size quantization energy of the dot). Therefore the dominant mechanism which leads to decoherence and is responsible for spin decoherence times being typically much smaller than spin relaxation times ($T_2 \ll T_1$) is the nuclear hyperfine interaction.

Research based on electronic transport through quantum dots was so far mainly focused on quantum dots fabricated on n-type GaAs heterostructures (for a review, see [33]). This is due to the high electronic quality and well developed processing technologies of n-type GaAs heterostructures. It has been demonstrated that the electron spin relaxation times in one-electron quantum dots can be as large as
Recent experiments [136] have also demonstrated the coherent manipulation of two coupled spins confined in a double quantum dot, and reported the electron decoherence time to be of the order \( T_2 \sim 10 \text{ ns} \). This value is limited by the nuclear hyperfine interaction between the confined electron spin and around \( 10^6 \) of the spin-3/2 nuclei of the host GaAs material.

In p-type GaAs heterostructures nuclear hyperfine interaction of holes with lattice nuclei is expected to be suppressed due to the p-like orbital symmetry of hole’s wave functions near each atom in the crystal and lack of overlap between hole and nuclear wavefunctions. Therefore, the decoherence of hole spins due to nuclear hyperfine interactions should be negligible. However, hole-doped GaAs low-dimensional systems have received little attention in efforts to utilize the spin in quantum information technologies. The reason was the well known fact that stronger spin-orbit interaction in hole-doped systems compared to their electron-doped counterparts significantly reduces spin relaxation times in bulk p-doped GaAs systems. Holes confined into quantum wells have much larger spin relaxation times (10-100 ps) compared to those in bulk p-type GaAs [137], but their spin relaxation still remains several orders of magnitude faster than electron spin relaxation. Recently, Bulaev and Loss predicted [138], that further confinement of holes into quantum dots can significantly increase the relaxation time \( T_1 \) of hole spins, so that it can be comparable, or even larger than that of electron spins. Besides, they have shown that even in the case of hole quantum dots the spin decoherence times due to spin-orbit coupling remains 2 times longer than spin relaxation time at low temperatures. In order to understand this increase of the spin relaxation times of carriers confined in quantum dots, one needs to keep in mind the fact that spin-orbit induced spin relaxation occurs actually through electron-phonon (or hole-phonon) interaction which affects spin via spin-orbit coupling. Due to the discrete energy spectrum of quantum dots, at very low temperatures \( k_B T \ll \Delta \), \( \Delta \) - the size quantization energy of the dot) electron (hole)-phonon scattering becomes much less effective than in the case of bulk or 2D structures, and therefore spin relaxation in quantum dots is much slower. It is important to mention that in order to keep the spin relaxation time long in the case of hole quantum dots, besides the condition \( k_B T \ll \Delta \), the lateral size of the dot \( l \) has to be much larger than the vertical size of the dot \( d \), \( l \gg d \), so that the mixing between heavy hole (HH) and light hole (LH) subbands is negligible, and that the dot has purely heavy hole character [138].

The manipulation of spins in quantum dots is typically realized with external magnetic fields produced by large magnetic coils. This conventional approach faces the problem of non-scalability. Therefore, there is an on-going investigation of the possibilities to manipulate spins in quantum dots in an all-electrical way. Spin-orbit interaction enables spins to be coupled to electric fields. Schemes for spin-orbit mediated, all-electrical coherent control of spins in quantum dots have recently been proposed [139, 140]. Hole quantum dots, due to their strong spin-orbit interaction, represent promising ground for exploring these concepts. However, one needs to keep in mind that holes confined in quantum dots are heavy holes with spin \( \pm 3/2 \),
and therefore the concepts developed for the manipulation of electron spins $\pm 1/2$ might be different in the case of holes due to the different selection rules in that case [141].

Hole quantum dots are relevant not only for the investigation of spin-orbit interactions, but also for studying carrier-carrier Coulomb interaction effects in low-dimensional systems. Due to the large effective mass of holes in GaAs, the importance of carrier-carrier interactions relative to the kinetic energy increases, and the interaction parameter $r_s$ becomes large, $r_s = E_{int}/E_F \gg 1$. It can therefore be expected, that the energy spectra of hole quantum dots can no longer be described by the constant interaction model [35] and that many-body effects must be involved. It is predicted that strong carrier-carrier interactions influence Coulomb blockade peak-spacing statistics in quantum dots [142]. Therefore studying energy spectra and peak-spacing fluctuations in hole quantum dots would provide new information about many-body effects in low-dimensional systems.

In this chapter we report the first Coulomb blockade (CB) measurements in single-hole transistors, defined on a p-type carbon doped GaAs heterostructure. The experimental investigation of electronic transport in p-type GaAs quantum dots was lacking so far due to technological difficulties to fabricate stable p-type structures with conventional split-gate techniques. Therefore, we have employed AFM oxidation lithography to fabricate p-type quantum dots. Using this fabrication technique significantly improved the stability of the fabricated p-type devices. Results from two quantum dots with different geometries, rectangular and ring-like, are presented. Both dots are fabricated on the same wafer (Bochum 12029), with the hole gas located 45 nm below the sample surface.

### 8.2 Rectangular quantum dot

We first study hole transport in a rectangular quantum dot. An AFM micrograph of the dot is shown in Fig. 8.1(a). The lithographic dimensions of the dot are $430 \times 170$ nm$^2$, while the width of both quantum point contacts (qpc) is $\sim 140$ nm. The height profile in Fig. 8.1(b) along the dashed line in (a) indicates excellent topological homogeneity and quite constant height of the oxide lines.

It is shown that for a 2DHG 45 nm below the sample surface AFM written oxide lines with a height of 15-18 nm completely deplete the 2DHG beneath at 4.2 K. For voltages in the range [-500 mV, +400 mV] applied to the plunger gate, the total leakage current to all other gates, which are kept grounded, is less than 1 pA (Fig. 8.1(c)). This allows us to make quantum dots with highly tunable in-plane gates.

#### 8.2.1 Coulomb blockade measurements

The transport measurements in the dot have been performed in a $^3$He/$^4$He dilution refrigerator at a base temperature of $\sim 50$ mK. We have measured the two-terminal
8.2. Rectangular quantum dot

Figure 8.1: (a) AFM micrograph of the quantum dot with designations of the gates: qpc 1 and qpc 2 are the gates for tuning the coupling of the dot to the source and drain, while the plunger gate serves to tune the number of holes in the dot. Bright oxide lines fabricated by AFM oxidation lithography lead to insulating barriers in the 2DHG. (b) Height profile along the dashed line in Fig. 1a - the oxide segments are 15-18 nm high. (c) Test of the insulating behavior of the oxide line at T=4.2 K: A voltage is applied to the plunger gate and the total current to all other gates, which are kept grounded, is measured. For this sample the oxide lines are insulating for the applied voltages in the range [-500 mV, +400 mV].

The conductance through the dot by applying either a small dc or ac bias voltage $V_{sd}$ between source and drain, and measuring the current through the dot with a resolution better than 50 fA.

The coupling of the dot to the source and drain contacts has been initially symmetrized by simultaneously sweeping qpc gates $V_{qpc1}$ and $V_{qpc2}$ at a fixed plunger gate voltage of $V_{pg} = -330$ mV (Fig. 8.2(a)). The symmetric configuration is achieved for $V_{qpc1} = -213$ mV, $V_{qpc2} = -236$ mV. It is also important to notice in figure 8.2(a) the change of the lever arms of the two qpc gates - this clearly shows that the dot is formed between the two qpc gates, and that the confinement is not dominated by one of the qpc gates.

The differential conductance of the dot is then measured as a function of the plunger gate voltage. Pronounced Coulomb resonances are observed (Fig. 8.2(b)). It is important to point out that the dot closes when the value of the plunger-gate voltage increases – this is a clear indication that we measure hole transport. In this configuration of the gate voltages the peak positions were stable in twenty consecutive plunger-gate sweeps within an accuracy of 0.1 mV. However, not every gate configuration shows such stability - in certain configurations charge rearrangements make reproducible measurements difficult, as it can be seen in Fig. 8.2(a).
Figure 8.2: (a) Differential conductance of the dot as a function of the two qpc gate voltages controlling the coupling to source and drain at the fixed plunger-gate voltage $V_{pg} = -330$ mV. (b) Differential conductance through the dot as a function of plunger gate voltage — clear Coulomb resonances are observed. Measurements are performed in the dot configuration: $V_{qpc1} = -213$ mV, $V_{qpc2} = -236$ mV, with symmetrically applied AC source-drain bias of 20 $\mu$V and frequency of 31 Hz at the base temperature of 50 mK. Inset: Blow up of the weak coupling regime.

8.2.2 Resonance peak shape and hole temperature

We now focus on the weak coupling regime shown in the inset of Fig. 8.2(b). Each of these five resonances is fitted both with an expression for a thermally broadened Coulomb blockade peak in the multi-level transport regime and a coupling broadened Lorentzian peak \[35\] (Fig. 8.3). In all cases the thermally broadened resonance fits significantly better to the data than a coupling broadened resonance, indicating that the dot is really in the weak coupling regime and that the peak broadening is determined by temperature rather than coupling. As a fit parameter we obtain FWHMs for each of these five peaks and they are in the range from 440 $\mu$V to 480 $\mu$V, which correspond to hole temperatures in the range $T_{\text{hole}} = 300 - 330$ mK. At this point it is important to address the question why the extracted hole temperature, which determines the peak broadening, is much higher than the mixing chamber temperature. It is already briefly mentioned in the previous paragraph, and will be discussed in more detail later, that low-frequency charge fluctuations are present in the transport measurements through the p-type dots. Besides these low-frequency fluctuations which can be seen in the transport measurements, it is possible that high-frequency charge fluctuations are also present. These fluctuations might be too fast to be resolved in the transport measurements, but they could contribute to the heating of holes. Another reason for the high hole temperature can be the fact that these measurements were performed in a dilution refrigerator with the insert still in a phase of optimization, and with the thermal coupling between the sample and the mixing chamber not fully optimized. Measurements performed in another
8.2. Rectangular quantum dot refrigerator, which will be presented later, indicate that the second reason is more probable.

\[ G_{\text{diff}} = \frac{2 \alpha e (V_{\text{pg}} - V_{\text{max}})}{2.5 k_B T_{\text{hole}}} \]

Figure 8.3: Peak at \( V_{\text{pg}} = -363.8 \text{ mV} \) (circles) fitted to a thermally broadened Coulomb blockade resonance (full line) and to a Lorentzian (dashed line). In the thermally broadened fit \( \alpha \) is the lever-arm of the plunger gate, while \( V_{\text{max}} \) and \( T_{\text{hole}} \) are fitting parameters.

8.2.3 Coulomb diamonds

Coulomb diamond measurements, i.e., measurements of the differential conductance as a function of bias voltage \( V_{\text{bias}} \) and plunger gate voltage \( V_{\text{pg}} \), are performed in the weak coupling regime, and the results are shown in Fig. 8.4(a). The uniform size of the diamonds indicates that all confined holes reside in one single potential minimum rather than occupying several disconnected or tunnel-coupled potential minima, as it was reported for a p-type SiGe quantum dot [143]. From the extent of the diamonds in bias direction we estimate a charging energy of the dot to be \( E_C \approx 1.5 \text{ meV} \), while the lever-arm of the plunger gate is \( \alpha \approx 0.26 \). This charging energy corresponds to a capacitance of the dot \( C = e^2/E_C \approx 1.1 \times 10^{-16} \text{ F} \). If we assume a disk-like shape of the dot, the capacitance is given by \( C = 8 \varepsilon_0 \varepsilon_r r \), where \( r \) is the radius of the dot. This allows us to estimate the electronic diameter of the dot to be \( \sim 230 \text{ nm} \), which is in very good agreement with the lithographic dimensions of the dot and indicates that the dot is really formed in the region encircled by the oxide lines. The estimated number of holes confined in the dot is \( \sim 150 \).

It can be seen in Fig. 8.4(b) that in certain dot configurations, low-frequency switching noise due to charge rearrangements in the sample becomes quite expressed. Although for low biases applied across the dot, the gate configurations can be found in which the switching noise in current through the dot can be completely suppressed, for large source-drain biases the switching noise becomes more pronounced.
Figure 8.4: (a) Coulomb diamonds in differential conductance, represented in a logarithmic gray scale plot (white regions represent low conductance). A DC bias is applied symmetrically across the dot and the current through the dot is measured. The differential conductance is calculated by numerical derivation. The charging energy of the dot is estimated to be $\sim 1.5$ meV from this measurement. Measurements are performed in the dot configuration: $V_{qpc1} = -225$ mV, $V_{qpc2} = -235$ mV. (b) Switching events in the dot current at higher biases for $V_{pg} = -340$ mV.

and is present in almost all dot configurations. The nature of the charge traps, responsible for the low-frequency switching noise is not fully understood. Most probably it is related to charging and discharging of the dopant acceptor levels. One can also imagine that disorder in the confining potential can produce such charge traps. In Si and SiGe quantum dots disorder can be very strong and therefore it can happen that charges reside in several tunnel-connected puddles instead of one potential minimum. However due to the high mobility of the GaAs heterostructure used for fabrication of the dot, this scenario seems less probable in our quantum dot.

In the case of the dot with steep potential walls, the mean single-particle level spacing can be calculated as $\Delta = \frac{2\pi \hbar^2}{g m^* A}$, where $g$ is the degeneracy of hole states and $A$ is the electronic area of the dot. Due to the large effective mass of holes ($m_1 = 0.34 m_e$ and $m_2 = 0.53 m_e$), the mean single-particle level spacing in the dot is estimated to be $\Delta \leq 15 \, \mu$eV, which is one order of magnitude smaller than typical values in electron quantum dots. Since we have $k_B T_{hole} \approx 25 \, \mu$eV for the estimated hole temperature, the dot is in the regime where $\Delta \leq k_B T_{hole}$. This explains why excited states cannot be resolved in Coulomb diamond measurements.

8.2.4 Temperature evolution of the Coulomb peaks

We further explored the temperature dependence of the Coulomb peaks and found that the peak amplitude does not decrease, but rather increases as the temperature increases (Fig. 8.5(a)). Such anomalous temperature behavior can be observed
when the temperature smearing of the Fermi-function in source and drain becomes comparable to a neighboring single-particle level spacing and additional channels enter the transport. This is another indication that the dot is not in the single-level transport regime, but in the intermediate multi-level transport regime where several (of the order of ten) single-particle levels participate in transport [33, 144].

8.2.5 Peak-spacing fluctuations

Despite the fact that the dot is in the multi-level transport regime, fluctuations in peak spacings are observed. The positions (in $V_{pg}$) of the five peaks shown in the inset of Fig. 2 are respectively: -369.6 mV, -363.8 mV, -357.7 mV, -353.3 mV, and -347.5 mV. Thus, the peak separations between consecutive peaks are: 5.8 mV, 6.1 mV, 4.4 mV and 5.8 mV, which after multiplying with lever-arm of the plunger gate gives the following separations in energy: 1.51 meV, 1.59 meV, 1.14 meV and 1.51 meV. These fluctuations of the peak spacings are much larger than the estimated single-particle level spacing and thus contradict the expectation of the constant interaction—random matrix theory. Similar observations have been made in a Si quantum dot [145] with interaction parameter $r_s = 2.1$ and similar charging energy and single-particle level spacing to the values in our device, and obtained results are attributed to electron-electron interactions. Since the interaction parameter in our dot is even larger, $r_s = 5$, the large fluctuations in peak spacings which we observe may indicate the significance of stronger carrier-carrier interactions in hole quantum dots.
8.2.6 Second cool-down

In order to understand the reasons for the relatively high hole temperature in our measurements, we performed a second cool-down of the same dot in another refrigerator where the thermal coupling of the sample to the mixing chamber was better and electronic temperatures around 100 mK were routinely obtained. Narrower Coulomb resonances were obtained (Fig 8.6) and the hole temperature extracted from the fitting of the peaks in this case was in the range $T_{\text{hole}} = 100 - 130$ mK. This indicates, that the peak width is really determined by temperature, and that there is no intrinsic mechanism in the sample which widens the Coulomb peaks up to this energy range [143]. Even though the hole temperature was lower in the second compared to the first cool-down, it was still not possible to resolve excited states in Coulomb diamond measurements. Therefore, in order to be able to investigate the single-particle level spectrum in hole quantum dots, one has to significantly reduce the lateral dimensions of the dot. It is also important to mention that the reduction of the hole temperature has not improved the stability of the dot.

![Figure 8.6: Coulomb blockade resonances in the second cool-down.](image)

Figure 8.6: Coulomb blockade resonances in the second cool-down. The dot is in the configuration $V_{qpc1} = -190.8$ mV, $V_{qpc2} = -178.8$ mV, with symmetrically applied DC source-drain bias of 10 $\mu$V; Inset: Fitting of the peak with thermally broadened Coulomb blockade resonance - extracted hole temperature is 100 mK

8.3 Coulomb blockade in ring-like quantum dot

The high tunability of our AFM defined devices is most strikingly demonstrated in the ring-like sample shown in the figure 8.7(a). This sample can be tuned by means of the in-plane gates from the completely open regime (four terminal resistance $\sim 20$ k$\Omega$) in which Aharonov-Bohm oscillations were studied (see chapter 9) to the completely closed, Coulomb blockaded regime, which will be analyzed in this
chapter. The lithographic radius of the ring is \( \sim 320 \text{ nm} \). The two-terminal ring conductance is measured in a dilution refrigerator with the base temperature \( \sim 70 \text{ mK} \). The coupling of the dot to source and drain is symmetrized and the qpc gates are set to the following values: \( V_3 = 72 \text{ mV}, V_4 = 120 \text{ mV}, V_5 = 310 \text{ mV}, V_6 = 200 \text{ mV}, V_{pg2} = -32 \text{ mV} \). Clear Coulomb oscillations are obtained upon sweeping plunger gate 1 (Fig 8.7(b)).

![Figure 8.7](image)

**Figure 8.7:** (a) AFM micrograph of the ring-like quantum dot with designation of the gates. (b) Coulomb blockade resonances in the configuration \( V_{pg2} = -32 \text{ mV}, V_3 = 72 \text{ mV}, V_4 = 120 \text{ mV}, V_5 = 310 \text{ mV}, V_6 = 200 \text{ mV}, \) applied DC bias 20 \( \mu \text{V} \). (c) Fitting of the peak with thermally broadened Coulomb blockade resonance (blue curve) and coupling broadened Lorentzian (red curve). (d) Coulomb diamonds in differential conductance for the ring-like dot in the configuration given in (b) — white regions represent low conductance. (e) Evolution of Coulomb blockade resonances in a perpendicular magnetic field.

Each of these resonances is fitted with a thermally-broadened Coulomb blockade peak and a coupling-broadened Lorentzian peak. However, in this case one can’t clearly distinguish which of the two functions fits better to the data (Fig 8.7(c)), which means that the coupling of the dot to the source and drain is not as weak as in the case of the rectangular dot. Therefore, the extracted hole temperature of \( T_{hole} = 160 \text{ mK} \), obtained from the fitting parameter FWHM = 425 \( \mu \text{V} \), represents the upper limit of the hole temperature, since the peak broadening is, besides temperature,
affected also by the coupling to the leads ($\Gamma = 195 \mu\text{V}$).

We have further measured nonlinear transport through the dot. From the obtained Coulomb diamonds (Fig. 8.7(d)) we estimate the charging energy $E_{C,\text{ring}} \approx 0.5$ meV and the lever arm of the plunger gate $1 \alpha_{\text{ring}} \approx 0.14$. From the value of the charging energy we estimate the electronic radius of the ring-like dot to be $r_{\text{ring}} \approx 340\text{nm}$, which is slightly larger compared to the lithographic radius of the ring.

Beside the central, white diamonds, where the current through the dot is blocked and the number of holes in the dot is a fixed integer number $N$, it is interesting to notice in the figure 8.7(d) well defined, light grey, side diamonds, where the current is enhanced due to the fact that the number of holes in the dot can fluctuate between $N$ and $N + 1$ in these regions of the stability diagram.

Excited states are not resolved either in these Coulomb diamond measurements. This is expected, since the mean single-particle level spacing, estimated from the electronic radius of the ring, is $\Delta_{\text{ring}} \leq 2 \mu\text{eV}$. Therefore, the ring-like dot is in the multi-level transport regime.

The evolution of the Coulomb peaks as a function of a perpendicular magnetic field is shown in Fig. 8.7(e). Although we clearly see the oscillations in the ring resistance with the Aharonov-Bohm period of $\Delta B \approx 60 \text{mT}$ in the open ring (see chapter 9), in the Coulomb blockade regime we do not see any signature of the Aharonov-Bohm effect and the peak positions do not shift with applied perpendicular magnetic field. This is presumably due to the fact that the dot is in the multi-level transport regime and the effect of the Aharonov-Bohm phase on different single-particle levels averages out.

Finally, we should again notice the presence of switching noise in the Coulomb diamond measurements in this dot (Fig. 8.7(d)), similar to the low-frequency noise observed in the rectangular dot. This is a clear indication that the traps which produce these switching events are not related to the particular device or some local defect, but are rather inherent to the wafer itself on which the structures are fabricated.

### 8.4 Summary and outlook

We have fabricated tunable quantum dots on a p-type GaAs/AlGaAs heterostructure by AFM oxidation lithography. By using this fabrication technique we were able to overcome the problems with large hysteresis effects present in gate sweeps in conventional split-gate defined nanostructures on p-type GaAs. The functionality of the dots is demonstrated by observing clear and reproducible Coulomb resonances in both dots, for the first time in p-type GaAs samples. The charging energies of the dots, determined from the Coulomb diamond measurements, are compatible with the lithographic dimensions of the dots.

Due to the large effective mass of the holes the single-particle level spacing in case
of our hole quantum dots is estimated to be smaller than the thermal broadening, and therefore it was not possible to resolve excited states in the differential conductance measurements. In order to be able to investigate the single-particle level spectrum in hole quantum dots, one has to significantly reduce both, the lateral dimensions of the dot as well as the hole temperature.

However, the fabrication of ultra-small hole quantum dot in which single-particle levels could be resolved remained as a challenge so far. One of the reasons is that for the required height $\sim 15$ nm of the oxide lines to be insulating, the width of the lines is typically $\sim 100$-$120$ nm and this is the limiting factor for the lateral dimensions of the dot. Combining the AFM defined in-plane gates with homogeneous metallic top gate for changing a global density of the structure, might be the way to reduce the size of the dot. Our tests showed that top-gates evaporated directly to the surface of the shallow p-type GaAs heterostructures suffer from current-leakage to the 2DHG and increase instability of the samples. However, putting additional insulating layer between the top-gate and the surface of the structure might improve characteristics of top-gated samples. This technological step requires further study and is not explored in this thesis.

Exploring single-particle level spectra in hole quantum dots would bring new information about the importance of carrier-carrier and spin-orbit interactions in low dimensional systems. Strong carrier-carrier Coulomb interactions are expected to greatly influence Coulomb peak spacings statistics [142]. In the presence of strong spin-orbit interactions different spin states of the dot can be mixed, and induced anticrossings in the level spectra can provide a measure of spin-orbit coupling strength, as was observed recently in InAs nanowire quantum dots [146].

The investigation of hole transport in a double quantum dot might provide the way to resolve single-particle levels in hole quantum dots, even when the level separation is smaller than the temperature broadening. In a double dot system single-particle levels can be detected through the transitions of holes from one dot to the other, rather than between a dot and a lead with thermally broadened density of states, as is the case for a single dot [147, 148]. Therefore, at large enough biases applied across the double dot, $eV_{\text{bias}} \gg k_B T$, transitions between the levels in the dots should be observable in a stability diagram of a double dot, even though the thermal broadening of the states in the leads might be larger than the level-separation in the dots.
Chapter 9

Aharonov-Bohm oscillations in the presence of strong spin-orbit interaction

9.1 Introduction - Spin-orbit induced geometric phases

Interference phenomena with particles have challenged physicists since the foundation of quantum mechanics. A charged particle, traversing a ring-like mesoscopic structure in the presence of an external magnetic flux $\Phi$, acquires a quantum mechanical phase due to the non-local interaction between the charge of the particle and the applied magnetic field with a vector potential $\vec{A}$. The interference phenomenon based on this phase is known as the Aharonov-Bohm (AB) effect \[7\], and manifests itself in oscillations in the resistance of the mesoscopic ring with periodicity $\Phi/\Phi_0$, where $\Phi_0 = h/e$ is a flux quantum. This Aharonov-Bohm phase is later seen as a special case of the geometric phase \[5, 6\] acquired by the orbital wave function of the charged particle encircling a magnetic flux line.

The particle’s spin can also acquire an additional geometric phase in systems with spin-orbit (SO) interaction \[15–20\]. The investigation of this spin-orbit induced phase in solid-state systems is currently the subject of intensive experimental work \[10–14, 149\]. The common point of these experiments is the investigation of electronic transport in ring-like structures defined on two-dimensional (2D) semiconducting systems with strong SO interaction. In such systems, an inhomogeneous, momentum dependent intrinsic magnetic field $\vec{B}_{\text{in}}$, perpendicular to the particle’s momentum, is present in the reference frame of the moving carrier \[43\]. The total magnetic field seen by the moving carrier is therefore $\vec{B}_{\text{tot}} = \vec{B}_{\text{ext}} + \vec{B}_{\text{in}}$, where $\vec{B}_{\text{ext}}$ is the external magnetic field perpendicular to the 2D system and $\vec{B}_{\text{in}}$ is intrinsic magnetic field in the plane of the 2D system. The particle’s spin precesses around $\vec{B}_{\text{tot}}$ and accumulates an additional phase upon cyclic evolution. This phase con-
9.2 Large ring

The sample was fabricated by AFM oxidation lithography on a p-type carbon doped (100) GaAs heterostructure, with a shallow 2DHG located 45 nm below the surface.
Chapter 9. Aharonov-Bohm oscillations in the presence of strong spin-orbit interaction

[150]. An AFM micrograph of the ring structure is shown in Fig. 9.1. The radius of the mean circular path within the ring is \( r = 420 \) nm, and the lithographic width of the arms is \( \sim 190 \) nm, which, taking into account the side depletion of the oxide lines, corresponds to an electronic width of the arms \( \sim 60 - 70 \) nm. The hole density in an unpatterned sample is \( 3.8 \times 10^{11} \) cm\(^{-2}\) and the mobility is 200 000 cm\(^2\)/Vs at a temperature of 60 mK. Therefore the Fermi wavelength is \( \sim 40 \) nm, and the mean free path is \( \sim 2 \) \( \mu \)m.

![AFM micrograph of the ring with designations of the gates. Bright oxide lines fabricated by AFM oxidation lithography lead to insulating barriers in the 2DHG.](image)

The presence of strong spin-orbit interactions in the heterostructure is demonstrated by a simultaneous observation of the beating in Shubnikov-de Haas oscillations and a weak anti-localization dip in the measured magnetoresistance of the Hall-bar fabricated on the same wafer (see chapters 5 and 6). The densities of the two spin-split subbands are deduced from the Fourier transform of Shubnikov-de Haas oscillations and they are \( N_1 = 1.35 \times 10^{11} \) cm\(^{-2}\) and \( N_2 = 2.45 \times 10^{11} \) cm\(^{-2}\). This allows us to quantify the strength of the Rashba spin-orbit interaction, assuming a cubic wave vector dependence for the splitting of heavy-hole band in (100) plane, \( \Delta_{SO} = 2\beta k_3^2 \) [43]. We obtain \( \beta = 2 \times 10^{-28} \) eVm\(^3\), which gives the spin-orbit induced splitting of the heavy hole subband at the Fermi level to be \( \Delta_{SO} \approx 0.8 \) meV. Due to the large effective mass of the holes \( (m_1 = 0.34 m_e \) and \( m_2 = 0.53 m_e) \) [65] the Fermi energy in the system, \( E_F = 2.5 \) meV, is much smaller compared to the Fermi energy in electron systems with the same density. Therefore the spin-orbit induced splitting of the heavy hole subband is \( \sim 30\% \) of the Fermi energy, documenting the presence of very strong spin-orbit interaction in the system.
9.2.1 Beating of Aharonov-Bohm oscillations

We have measured the four-terminal resistance of the ring in a $^3$He/$^4$He dilution refrigerator at a base temperature of $\sim 60$ mK with lock-in technique. A low ac current of 2 nA and frequency of 31 Hz was applied through the sample, in order to prevent sample heating. We found that holes in GaAs are extremely sensitive to heating and that an increase of the current from 2 nA to 10 nA reduces the amplitude of the observed oscillations by more than a factor of 5. Since the corresponding voltage oscillations were less than 100 nV at the current level of 2 nA we mounted a low-noise voltage amplifier with a gain of 1000 directly at the top of the fridge in order to suppress the noise from the cabling as much as possible.

Figure 9.2: (a) Measured magnetoresistance of the ring (blue curve) together with the low-frequency background resistance (red curve); (b) AB oscillations obtained after subtraction of the low-frequency background from the raw data. A clear beating pattern is revealed in the AB oscillations. (c) Fourier transform spectra of the AB oscillations revealing $h/e$ and $h/2e$ peaks. (d) Splitting of the $h/e$ Fourier peak. (e) Splitting of the $h/2e$ Fourier peak
Fig. 9.2(a) shows the measured magnetoresistance of the ring (blue curve) together with a low-frequency background resistance composed of the low-frequency Fourier components of the signal (red curve). Aharonov-Bohm (AB) oscillations are clearly resolved, their peak-to-peak amplitude being \( \sim 200 \ \Omega \) on a background of \( \sim 6 \ \text{k}\Omega \). This corresponds to a visibility larger than 3%. We restrict the measurements of the AB oscillations to magnetic fields in the range from \((-0.2 \ \text{T to } 0.2 \ \text{T})\) in order to prevent their mixing with SdH oscillations, which start to develop above 0.2 T at a temperature of 60 mK. Before the measurement, qpc gates 3, 4, 5 and 6 are set to fixed voltages, so that the background resistance is \( \sim 6 \ \text{k}\Omega \), and they are kept at the same values throughout all the measurements presented in this chapter. Plunger gates \( pg_1 \) and \( pg_2 \) are set to \( V_{pg_1} = -145 \ \text{mV} \) and \( V_{pg_2} = -95 \ \text{mV} \) in the measurements presented in Fig. 9.2(a), and their influence on the AB oscillations will be discussed later.

After subtracting the low-frequency background from the raw data a clear beating pattern is revealed in the AB oscillations with a well defined node at \( \sim 115 \ \text{mT} \) (Fig. 9.2(b)), where a phase jump of \( \pi \) occurs (arrow in Fig. 9.4(c)). The position of the beating node indicates the presence of two oscillation frequencies differing by \( 1/0.115 \approx 9 \ \text{T}^{-1} \). The Fourier spectrum of the AB oscillations, taken in the symmetric magnetic field range \((-0.2 \ \text{T}, 0.2 \ \text{T})\), reveals an \( h/e \) peak around 130 T\(^{-1}\) (Fig. 9.2(c)), corresponding to a period of the AB oscillations of 7.7 mT and the radius of the holes’ orbit of 415 nm, which is in excellent agreement with the lithographic mean radius of the ring.

### 9.2.2 Dependence of the Fourier spectra of the Aharonov-Bohm oscillations on magnetic field range

If we zoom in on the \( h/e \) peak (Fig. 9.2(d)) in the Fourier spectrum, the splitting of this peak into 3 peaks at the frequencies 127 T\(^{-1}\), 136 T\(^{-1}\) and 143 T\(^{-1}\) can be observed. We have carefully checked that the splitting of the Fourier peak is genuine to the experimental data and not a result of the finite data range, by reproducing the same splitting with different window functions for the Fourier transform. The differences of the oscillation frequencies agree with that anticipated from the position of the beating node in the raw data.

Similar splittings in the \( h/e \) peak observed in [11] were attributed to the signature of the spin Berry phase, although in that case the beating in the raw-data was not clearly seen, due to a small amplitude of the AB oscillations. In reference [12] it was discussed that a splitting of the \( h/e \) peak in the Fourier spectrum can also be obtained without taking into account the spin Berry phase and therefore the reliability of a Fourier transform, without the observed beating in the raw data, is limited. We explored in detail the dependence of the Fourier spectrum on the magnetic field range in which the Fourier transform is performed for two different gate configurations of the ring (Fig 9.3). Upon reducing the symmetric range of magnetic fields in which the Fourier transform is performed, the positions of the peaks remain
fixed, while their amplitudes become smaller (Fig 9.3(b) and 9.3(e)). On the other hand keeping the B-range in which the Fourier transform is performed constant and shifting it along the B-axis produces more complex changes in Fourier transforms (Fig 9.3(c) and 9.3(f)). The observed complex dependence of AB Fourier spectrum on the magnetic field range in which the Fourier transform is performed is similar to what has been reported in [11]. It can be understood from the fact that the direction of the effective magnetic field $\overrightarrow{B}_{eff}$, around which the spin precesses, continuously changes upon changing external magnetic field, and therefore the geometric phase acquired by the spins will be different for different magnetic field ranges in which analysis is performed.

We also observed that changes in the relative phase between the two arms of the ring induced by changes of plunger-gate voltages significantly influence the splitting of the h/e peak. Although in most gate configurations the triple-peak pattern around

Figure 9.3: Left column corresponds to the gate configuration $V_{pg1} = -145$ mV and $V_{pg2} = -95$ mV: (a) Magnetoresistance after subtracting the low-frequency background; (b) Fourier transforms for symmetric B-field ranges, as indicated; (c) Fourier transforms for shifting constant B-field interval, as indicated. Right column corresponds to the gate configuration $V_{pg1} = -172$ mV and $V_{pg2} = -188$ mV: (d) Magnetoresistance after subtracting low-frequency background; (e) Fourier transforms for symmetric B-field ranges, as indicated; (f) Fourier transforms for shifting constant B-field interval, as indicated.
the h/e peak was visible, the relative strength of the peaks was strongly dependent on the plunger gate voltages (Fig 9.3(b) and 9.3(e)).

Therefore, rather than analyzing Fourier transforms [10, 11] whose details depend on the magnetic field range in which the Fourier transform is taken, we now focus directly on the magnetic field-dependent resistance.

### 9.2.3 Altshuler-Aronov-Spivak h/2e oscillations

In contrast to the h/e-periodic AB oscillations which are very sensitive to phase changes in the ring arms, Altshuler-Aronov-Spivak (AAS) h/2e oscillations originating from the interference of time reversed paths are expected to be more robust if the microscopic configuration of the arms is changed. However, h/2e oscillations have not been reported so far for GaAs holes. In Fig. 9.2(c) we can resolve a the peak at \( \sim 270 \, \text{T} \) in the Fourier spectrum, corresponding to h/2e oscillations. When zooming in on it (Fig. 9.2(e)) we see a clear splitting of this peak.

In Fig. 9.4(a) we present the raw data after subtracting the low-frequency background (red line) together with the filtered h/e oscillations (dashed line) [151]. The h/e contribution to the signal is the inverse Fourier transform of the h/e peak in the Fourier spectrum. In the further text we will use the following notation: \( R_d \) corresponds to the raw data, \( R_b \) is the low-frequency background, \( R_{h/e} \) is the inverse Fourier transform of the h/e peak and \( R_{h/2e} \) is the inverse Fourier transform of the h/2e peak in the Fourier spectrum. One can see that the raw data contains additional resistance modulations, beyond the h/e oscillations. In order to demonstrate that those additional features are due to the h/2e oscillations we plot in Fig. 9.4(b) the difference \( R_d - R_b - R_{h/e} \) (red line) and the curve \( R_{h/2e} \) obtained by inverse Fourier transform of the h/2e peak (dashed line) and find excellent agreement.

We further plot in Fig. 9.4(d) the difference \( R_d - R_b - R_{h/e} \) (red line), together with the filtered h/2e oscillations \( R_{h/2e} \) (dashed line) in a larger range of magnetic fields. A beating in the h/2e oscillations is observed, with nodes developing around 40 mT, 115 mT, and 165 mT (arrows). This kind of aperiodic modulation of the envelope function of the h/2e oscillations, rather than the regular beating, is predicted for the case of diffusive rings in the presence of Berry’s phase [15], since Berry’s phase also changes with increasing external magnetic field. Besides, the observed splitting of the h/2e peak in the Fourier spectrum with side peaks at 266 T\(^{-1}\) and 274 T\(^{-1}\) (Fig. 9.2(e)) also agrees qualitatively with the predictions made in [15] and points to the signature of the spin Berry phase in our ring. We point out that the separation between the split h/2e peaks of 8 T\(^{-1}\) is the same as the separation between the split h/e peaks. In the plot of the filtered h/e oscillations (Fig. 9.4(c)) we notice that only the node around 115 mT is common for both, the h/e and h/2e oscillations, while the other two nodes in the h/2e oscillations correspond to maxima in the beating of the h/e oscillations.
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Figure 9.4: (a) Measured magnetoresistance of the ring after subtracting the low-frequency background $R_d - R_b$ (full line, red) together with the filtered h/e oscillations $R_{h/e}$ (dashed line). (b) Difference $R_d - R_b - R_{h/e}$ (full line, red) together with the inverse Fourier transform of the h/2e peak $R_{h/2e}$ (dashed line). (c) Beating in the filtered h/e oscillations. The width of the gray and white rectangles corresponds to the period of 7.7 mT. The arrow points to the beating node where a phase jump of $\pi$ occurs. (d) Beating in the filtered h/2e oscillations with arrows indicating possible nodes.

9.2.4 Gate dependence of Aharonov-Bohm oscillations - Phase jumps

The evolution of the AB oscillations upon changing plunger gate voltages $V_{pg1}$ and $V_{pg2}$ is further explored (Fig. 9.5(a)). Plunger gate voltages are changed antisymmetrically: $V_{pg1} = -120 \text{mV} -V$; $V_{pg2} = -120 \text{mV} +V$. Two distinct features are visible: there is always a local minimum in the AB oscillations at $B = 0$ T, and the oscillations experience a phase jump by $\pi$ around $V = 27 \text{mV}$. In order to understand the origin of these two features we plot the filtered h/e (Fig. 9.5(c)) and h/2e oscillations (Fig. 9.5(d)) as a function of $V$. It can be seen that the h/e oscillations experience a phase jump of $\pi$ (Fig. 9.5(c)), while the h/2e oscillations do not (Fig. 9.5(d)). We have explored this behavior in another gate configuration.
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Figure 9.5: (a) Evolution of the AB oscillations upon changing plunger gate voltages $V_{pg1} = -120 \text{mV} - V$; $V_{pg2} = -120 \text{mV} + V$. (b) Examples of AB oscillations at $V = 25 \text{mV}$ (black line) and $V = 35 \text{mV}$ (red line) corresponding to dashed lines in (a). (c) Filtered $h/e$ oscillations as a function of plunger gate voltages, showing the phase jump of $\pi$ around $V = 27 \text{mV}$. (d) Filtered $h/2e$ oscillations as a function of plunger gate voltages, showing the local minimum at $B = 0 \text{T}$ at all gate voltages.

$V_{pg1} = -180 \text{mV} - V$; $V_{pg2} = -180 \text{mV} + V$ and found the same results (Fig. 9.6). If the plunger gate voltages $V_{pg1}$ and $V_{pg2}$ are changed symmetrically ($V_{pg1} = -120 \text{mV} + V$, $V_{pg2} = -120 \text{mV} + V$, or $V_{pg1} = -180 \text{mV} + V$, $V_{pg2} = -180 \text{mV} + V$) we see phase jumps neither in $h/e$ nor in $h/2e$ oscillations. The reason for such a behavior is the fact that the $h/e$ oscillations are sensitive to the phase difference $\Delta \varphi = k_1l_1 - k_2l_2$ between the two arms, which can be changed by plunger gates, while the $h/2e$ oscillations originate from the interference of exactly the same, time reversed paths which always interfere constructively, if spin-orbit interaction is neglected.

However, this constructive interference of partial waves propagating along the same paths in opposite directions should produce a resistance maximum in the $h/2e$ oscillations at $B = 0 \text{T}$, as it is the case for electrons in GaAs ring structures [152, 153]. The fact that we observe a resistance minimum at $B = 0 \text{T}$ in all gate configurations of the GaAs hole ring (Fig. 9.5(a)), which is due to a minimum at $B = 0 \text{T}$ in $h/2e$ oscillations (Fig. 9.5(d)), is a clear indication that the phase acquired by the spin part of the hole wave functions propagating in opposite directions plays a significant role in our system. This effect has exactly the same origin as the weak...
9.2. Large ring

Figure 9.6: (a) Evolution of the AB oscillations upon changing plunger gate voltages $V_{pg1} = -180 \text{mV} - V; V_{pg2} = -180 \text{mV} + V$. (b) Filtered h/e oscillations as a function of plunger gate voltages, showing the phase jumps of $\pi$. (c) Filtered h/2e oscillations as a function of plunger gate voltages, showing the local minimum at $B = 0 \text{T}$ at all gate voltages.

... anti-localization effect in the systems with strong spin-orbit interactions, and can be attributed to the destructive interference of spins propagating along time reversed paths [41]. It should be emphasized that this minimum is not caused by weak-antilocalization in the ring leads, since the weak-antilocalization dip in 2D samples has a much smaller magnitude (less than 1$\Omega$, see chapter 6) than the minimum at $B = 0 \text{T}$ in the ring.

Figure 9.7: Evolution of the (a) h/e Fourier peak, (b) h/2e Fourier peak, upon changing plunger gate voltages $V_{pg1} = -120 \text{mV} - V; V_{pg2} = -120 \text{mV} + V$. White in the color scale corresponds to low and black to high FFT amplitude. However, the absolute scales for the figures (a) and (b) are different.

As it was mentioned before the splitting of the h/e Fourier peak shows very complex behavior upon changing plunger gate voltages in the range shown in Fig. 9.5(a). In Fig. 9.7(a) we show a detailed dependence of the splitting of the h/e peak upon changing plunger gate voltages $V_{pg1} = -120 \text{mV} - V; V_{pg2} = -120 \text{mV} + V$. It can be seen that the relative strength of the split peaks changes upon crossing the phase jump around $V = 27 \text{mV}$, and that the triple-peak structure present below...
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$V = 27$ mV evolves in double or single peak structure above $V = 27$ mV. The h/2e peak shows a double peak structure, with peaks at $266$ T$^{-1}$ and $274$ T$^{-1}$ in most, but not all gate configurations (Fig. 9.7(b)).

![Graph showing magnetoresistance and Fourier transforms](image)

Figure 9.8: (a) Ensemble averaged magnetoresistance of the ring after subtracting the low-frequency background (red line) together with the filtered h/e (black dashed line) and h/2e component (blue line). (b) Splitting of the h/2e peak of the ensemble averaged data. (c) Beating in the h/2e component of the ensemble averaged data.

In order to check how robust the splitting of the h/2e Fourier peak is, we performed ensemble averaging over the plunger gate voltages in the range shown in Fig. 9.5(a). This procedure is justified, because we found that the ring background resistance does not change significantly (less than 10%) upon changing plunger gate voltages in this range. This means that the plunger gates do not influence significantly the electronic width of the arms (since they are too open), but only affect the phase picked up by the holes upon traversing a given ring arm. Therefore we can consider that upon changing the plunger gates in the range shown in Fig. 9.5(a) the ring remains in the same macroscopic state, and only its microscopic configurations change. After ensemble averaging the contribution from the h/2e oscillations becomes more visible (see Fig. 9.7(a)), due to partial (but not complete) suppression of the h/e oscillations. The Fourier transform of the ensemble averaged data also shows the splitting of the h/2e peak, with the side peaks at the same positions $266$ T$^{-1}$ and $274$ T$^{-1}$ (Fig. 9.7(b)). Signatures of a beating in the h/2e oscillations can also be seen in the ensemble averaged data (Fig. 9.7(c)).
Finally we have explored the temperature dependence of the AB oscillations in several different gate configurations and found that the oscillations persist up to 350 mK. The temperature evolution of the AB oscillations in the gate configuration $V_{pg1} = -155$ mV and $V_{pg2} = -85$ mV is shown in Fig 9.8(a).

The amplitude of the h/e oscillations at a given temperature is obtained from the numerical integration of the h/e Fourier peak in the frequency window $[0.5h/e, 1.5h/e]$. In Fig. 9.8(b) we plot the natural logarithm of such a h/e Fourier amplitude as a function of temperature. The relation between the amplitude of the h/e AB oscillations and phase coherence length $L_\phi$ is $R = R_0 \exp(-L/L_\phi)$, where $L$ is the characteristic length of the system and in the case of the ring structure is $L = 2r\pi$. If we assume that the phase-coherence length of the holes follows a $1/T$ temperature dependence, reported for the case of open electron AB rings [153], and theoretically predicted for the case of ballistic, one-channel rings [154], we can extract the phase coherence length of the holes to be $L_\phi = 1.3 \times 10^{-7}$[m·K]/$T$, from the linear fit of the data in Fig. 9.8(b). This gives $L_\phi = 2$ μm at the base temperature of $T = 65$ mK. We performed the same analysis in several other plunger-gate configurations of the ring and found that the values for the phase coherence length at $T = 65$ mK are in the range $L_\phi = 2.0 \pm 0.4$ μm. This value is approximately one order of magnitude smaller than the value of $L_\phi$ in electron AB rings with comparable densities and mobilities [152, 153]. Such a tendency was also observed in the recent measurements of dephasing times of holes in open quantum dots [155], and suggests stronger dephasing in hole compared to electron systems, presumably due to
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stronger carrier-carrier interactions in hole systems [154].

9.2.6 Aharonov-Bohm oscillations in the Quantum Hall regime

In this paragraph we present the measurements of the Aharonov-Bohm oscillations at higher magnetic fields in the Quantum Hall regime. Figure 9.10(b) shows the four-terminal magnetoresistance of the ring with well developed minima at filling factors 2 and 3. The fact that at filling factor 2 the resistance at the minimum goes to zero demonstrates that both edge states are transmitted through the ring. The resistance minimum at filling factor 3 does not go completely to zero, but it also remains far below the value of $\frac{h}{6e^2}$ which would correspond to one edge channel reflected back. This suggests that at filling factor 3 there might be an edge channel which is partially, but not fully reflected back from the ring. It should be mentioned that all measurements presented in this paragraph are performed in the following gate voltage configuration: $V_5 = V_6 = -180$ mV, $V_3 = V_4 = -300$ mV, $V_{pg1} = -300$ mV, $V_{pg2} = -3260$ mV, and the applied current was 2 nA.

![Figure 9.10: (a) AFM micrograph of the quantum ring. (b) Four-terminal magnetoresistance of the ring with indications of SdH minima at filling factors 2 and 3. Applied current was 2nA and the ring in-plane gates were in the following configuration: $V_5 = V_6 = -180$ mV, $V_3 = V_4 = -300$ mV, $V_{pg1} = -300$ mV, $V_{pg2} = -260$ mV.](image)

The ring magnetoresistance displays oscillatory behavior around the minimum at filling factor 2 (Fig. 9.11(a)). Zooming in the B-field regions (7.1T, 7.2T) and (7.2T, 7.3T) (Fig. 9.11(b),(d)) reveals well resolved AB oscillations. The Fourier transform shows in both cases a peak at 260 T$^{-1}$, which is exactly double the frequency observed in the low-field AB oscillations.

Similarly, the ring magnetoresistance around filling factor 3 (Fig. 9.12) shows well resolved AB oscillations with a frequency of 390 T$^{-1}$ (Fig. 9.12(c)), which is exactly three times larger than the frequency of the low-field AB oscillations. Therefore we can conclude that we observe the AB oscillations with a frequency $\frac{h}{2e}$ at filling factor 2, and with a frequency $\frac{h}{3e}$ at filling factor 3.

We have observed that outside of the minima in SdH oscillations, the AB oscillations become suppressed (see the left-hand side of Fig 9.11(a) and the right hand
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Figure 9.11: (a) Magnetoresistance of the ring around filling factor 2. (b) Aharonov-Bohm oscillations in the B-field region (7.1T, 7.2T). (c) Fourier transform of the AB oscillations shown in (b) with a peak around 260 T$^{-1}$. (d) AB oscillations in the B-field region (7.2T, 7.3T). (e) Fourier transform of the AB oscillations shown in (d) with a peak around 260 T$^{-1}$.

We did not find any B-field region between the minima at filling factors 2 and 3 where both $h/2e$ and $h/3e$ oscillations are simultaneously present. We found that changing of voltages applied to the qpc or arm gates of the ring does not influence the frequencies of the observed AB oscillations at filling factors 2 and 3, but only slightly shifts the positions of the SdH minima in magnetic field, where the oscillations are observed.

The observation of the double-frequency AB oscillations was first reported for quantum antidots in the quantum Hall regime, and it was explained in terms of charging of edge states circulating around an antidot [156]. This frequency doubling was later commonly observed for antidots in the quantum Hall regime. $h/3e$ oscillations have also been seen in AB measurements on antidot samples in the quantum Hall regime, when 3 edge channels encircle the antidot [157]. The magnetoresistance of a disk-like cavity in a 2DEG was studied in the Quantum Hall regime and


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4.3

1.5

1.0

0.5

0.0

5.0

4.5

4.0

3.5

3.0

2.5

2.0

1.5

1.0

0.5

0.0

5.0

4.5

4.0

3.5

3.0

2.5

2.0

1.5

1.0

0.5

0.0

$R (k \Omega)$

$B (T)$

Figure 9.12: (a) Magnetoresistance of the ring around filling factor 3. (b) AB oscillations in the B-field region (4.35T, 4.45T). (c) Fourier transform of the AB oscillations shown in (b) with a peak around 390 T$^{-1}$.  


different periods were observed at different filling factors $i$, satisfying approximately $i \cdot \Delta B = \Phi_0 / A$, where $A$ is the area of the ring and $\Phi_0 = \hbar/e$ [158, 159]. The observed Aharonov-Bohm-type oscillations in this disk-like device were attributed to the formation of the zero-dimensional states, formed due to interference of the edge channels along the circumference of the disk.

9.3 Small ring

We further explore Aharonov-Bohm oscillations in the small quantum ring (Fig. 9.13(a)) with the lithographic radius of 320 nm and the radius of the mean circular orbit in the ring of $\sim 160$ nm. The sample is fabricated on the same wafer (Bochum 12029) as the large ring sample, presented in the first part of this chapter, and therefore has about the same spin-orbit interaction strength, density and mobility. The sample is highly tunable and its resistance can be changed from $\sim 20$ k$\Omega$ to complete pinch-off by means of the in-plane gates. While in chapter 8 we investigated the same sample in the Coulomb blockade regime, here we focus on the more open regime and keep the four-terminal ring resistance in the range $20 - 50$ k$\Omega$. Before the measurements the qpc gates were set to the following values $V_3 = -100$ mV, $V_4 = -215$ mV, $V_5 = -360$ mV, $V_6 = -140$ mV, so that the ring is as open as possible, and these values were kept fixed throughout all measurements presented in this chapter. Further opening of the ring, by going to even more negative values for qpc gate voltages, was not possible due to the leakage of the oxide lines at lower qpc
9.3. Small ring

gate voltages. Measurements of the four-terminal ring resistance were performed in a $^3$He/$^4$He dilution refrigerator at a base temperature of $\sim 65$ mK with lock-in technique. In order to prevent sample heating we applied a low ac current of 0.5 nA and frequency of 31 Hz.

Figure 9.13: (a) AFM micrograph of the small ring with designation of the gates. (b) Measured magnetoresistance of the small ring (blue curve) together with low-frequency background resistance (red curve) in the plunger gate configuration $V_{pg1} = -95$ mV, $V_{pg2} = -150$ mV. (c) Fourier transform of the data with a strong $h/e$ peak at 17.5 T$^{-1}$ and weaker $h/2e$ and $h/3e$ peaks. (d) AB oscillations obtained after subtraction of the low-frequency background from the raw data (red line) together with the filtered $h/e$ oscillation (dashed line) (e) Difference between the data and filtered $h/e$ oscillations (red line) together with the curve obtained by inverse Fourier transform of the $h/2e$ peak (black dashed line).

Fig. 9.13(b) shows the measured magnetoresistance of the ring (blue line) together with a low frequency background resistance (red line) in the plunger gate configuration $V_{pg1} = -95$ mV, $V_{pg2} = -150$ mV. AB oscillations with a peak-to-peak amplitude of $\sim 4$ k$\Omega$ are clearly resolved (Fig. 9.13(d)). Therefore the visibility of the AB oscillations is larger than 10%, which is the highest value reported to date for holes [149]. The Fourier transform of the AB oscillations shows a strong $h/e$ peak at 17.5 T$^{-1}$ together with faint, but visible $h/2e$ and $h/3e$ peaks. Therefore the pe-
period of the AB oscillations is 57 mT, which corresponds to a radius of the holes’ orbit of 150 nm. This value is in very good agreement with a lithographic mean radius of the holes’ orbit within the ring. If we plot the filtered h/e oscillations, obtained by the inverse Fourier transform of the h/e peak in the Fourier spectrum [151] (dashed line in 9.13(d)) it is obvious that the data also contains contributions from higher harmonics. Fig. 9.13(e) shows the difference between the data and filtered h/e oscillations (red line) together with the curve obtained by the inverse Fourier transform of the h/2e peak (black dashed line). Good agreement between these two curves provides evidence for the presence of h/2e contribution in the magnetoresistance of the small ring.

In order to further support our claims we show the results from another plunger gate configuration $V_{pg1} = -78.5 \text{ mV}$, $V_{pg2} = -222 \text{ mV}$ in Fig. 9.14 and perform the same analysis as above. We observe even stronger h/2e and h/3e peaks in the Fourier spectrum in this compared to the previous gate configuration. The h/e peak exhibits a splitting into two equally strong peaks at 12 T$^{-1}$ and 17 T$^{-1}$. However, due to the very large period of the AB oscillation, only up to 10 oscillations are present in the magnetic field range (-0.3 T, +0.3 T) where the measurement is performed and no beating can be seen in the raw data. Therefore, although the amplitude of the AB oscillations in the case of the small ring is quite large, the fact that due to their larger period only few oscillations are present in the relevant B-field range, prevents us from performing a detailed analysis of the beating of the AB oscillations as we have done in the case of the large ring.

We analyze in detail the dependence of the AB oscillations on plunger gate configurations. Fig. 9.15(a) shows the four-terminal resistance of the ring as a function of plunger gate voltages $V_{pg1}$ and $V_{pg2}$ at $B = 0 \text{T}$ and fixed qpc gate voltages given above in the text. Due to the small size of the ring, the influence of the plunger gates on the ring resistance is much stronger than it was the case for the large ring sample, and the resistance of the small ring is tuned in the range 20 - 50 kΩ in the measurements presented in this chapter. The dependence of the AB oscillations on plunger gate configurations is investigated along the two lines in parameter space - along the line $V_{pg1} = 0.5 \cdot V_{pg2} - 20 \text{ mV}$ both plunger gate voltages increase and the resistance of the ring changes continuously from 20 - 50 kΩ, while along the other line $V_{pg1} = -0.75 \cdot V_{pg2} - 245 \text{ mV}$ an increase of one plunger gate voltage is followed by a decrease of the other plunger gate voltage and the resistance of the ring does not change significantly and remains around 30 kΩ. Gate changes along the first line induce a narrowing of the both ring arms, and effectively shrink the holes’ orbit inside the ring. The changes along the second line induce a narrowing of one ring arm and widening of the other, and therefore affect the symmetry of the ring, while keeping the size of the holes’ orbit inside the ring approximately unchanged.

The left column in Fig. 9.15 corresponds to the gate changes along the line $V_{pg1} = 0.5 \cdot V_{pg2} - 20 \text{ mV}$, and shows the evolution of the raw AB oscillations (b), filtered h/e component of the oscillations (c) and the difference between the data and the filtered h/e oscillations (d), respectively. Similarly, the right column of Fig.
Figure 9.14: (a) Measured magnetoresistance of the small ring (blue curve) together with low-frequency background resistance (red curve) in the plunger gate configuration $V_{pg1} = -78.5$ mV, $V_{pg2} = -222$ mV. (b) AB oscillations obtained after subtraction of the low-frequency background from the raw data (red line) together with the filtered $h/e$ oscillation (dashed line) (c) Fourier transform of the data with split $h/e$ peak and well defined $h/2e$ and $h/3e$ peaks.

9.15 shows the same dependencies along the line $V_{pg1} = -0.75 \cdot V_{pg2} - 245$ mV in parameter space.

The common point for both series of measurements is the presence of a local resistance minimum at $B = 0$ T in all gate configurations (Fig 9.15 (b) and (e)), which is related to the minimum in the $h/2e$ oscillations at $B = 0$ T (Fig 9.15 (d) and (g)). This behavior is exactly the same as observed in the large ring sample and further supports our claim that in the presence of strong spin-orbit interactions the hole spins propagating along the time reversed paths interfere destructively.

We again observe phase jumps for the $h/e$ oscillations (Fig 9.15 (c)) and no phase jumps for $h/2e$ oscillations (Fig 9.15 (d) and (g)) which is exactly the same behavior as observed in the large ring sample. However, in contrast to the case of the large ring where we observe phase jumps for $h/e$ oscillations for the antisymmetric changes of the plunger gates, and no phase jumps for the symmetric changes, in the case of the small ring sample we observe clear phase jumps for symmetric changes of the plunger gates (Fig 9.15 (c)), while an asymmetric change of the plunger gates does not produce such clear features (Fig 9.15 (f)). This behavior is counterintuitive, since phase jumps are more expected in the case of larger asymmetry between the ring arms.

The fact that the phase of the AB oscillations cannot change continuously, but only in discrete steps of $\pi$ is a consequence of the Onsager relations $G_{ij}(B) =$
Figure 9.15: (a) Resistance of the ring as a function of plunger gate voltages $V_{pg1}$ and $V_{pg2}$ at $B = 0$ T and fixed qpc gate voltages given in the text. Indicated are the two lines in parameter space along which the evolution of AB oscillations is explored. The left column shows measurements along the line $V_{pg1} = 0.5 \cdot V_{pg2} - 20 \text{mV}$: (b) Evolution of the AB oscillations upon changing plunger gate voltages, (c) Filtered h/e oscillations as a function of plunger gate voltages, (d) Difference between the data and filtered h/e oscillations as a function of plunger gate voltages. The right column shows measurements along the line $V_{pg1} = -0.75 \cdot V_{pg2} - 245 \text{mV}$: (e) Evolution of the AB oscillations upon changing plunger gate voltages, (f) Filtered h/e oscillations as a function of plunger gate voltages, (g) Difference between the data and filtered h/e oscillations as a function of plunger gate voltages.
9.4 Summary

We have measured highly visible Aharonov-Bohm oscillation in two quantum rings fabricated by AFM oxidation lithography on p-type GaAs heterostructure with strong spin-orbit interaction. The visibility of the AB oscillations in the large ring sample with a mean radius of 420 nm is about 3%, while in the case of the small ring with the radius of the orbit of 160 nm the visibility of the AB oscillations is larger than 10%, which is the highest value reported to date.

Beside the AB phase, acquired by the orbital part of the hole’s wave function, in a system with strong SO interaction the spin part of the hole’s wave function acquires an additional geometric phase. The interplay of these two phases is predicted to produce a complex beating-like oscillation pattern [15]. In contrast to previous experiments on p-type GaAs rings, where the signature of the phase acquired by the hole’s spin was attributed to the splitting of the h/e peak in the Fourier spectrum, we directly observe beating patterns, with clearly developed nodes, in the measured magnetoresistance of the large ring sample. Beside the h/e oscillations, we resolve the contribution from the h/2e oscillations and show that they also produce a beating pattern, resulting in the splitting of the h/2e Fourier peak. In the small ring sample we have also resolved h/2e and h/3e harmonics of the AB oscillations. However, due to the large period of the AB oscillations in the small ring, only few oscillations are present in the relevant B-field range which prevents us from analyzing a beating in the case of this sample.

An additional signature of the spin interference in our rings is the observation of a local minimum in the resistance at B=0T in all gate configurations of both, the large and the small ring. We show that it is related to a minimum in the h/2e oscillations.
Chapter 9. Aharonov-Bohm oscillations in the presence of strong spin-orbit interaction

Upon changing voltages applied to the arm-gates of the rings, we observe in both samples that the h/e oscillations experience phase jumps, while the h/2e oscillations do not and always show a minimum at $B = 0$T. This resistance minimum in the h/2e oscillations has the same origin as the weak anti-localization effect. The observed minimum is a result of the holes’ interference in the ring and it shows that upon one full turn along the time reversed paths the holes’ spins propagating in different directions acquire a phase difference of $\pi$ and interfere destructively.

The temperature dependence of the AB oscillations reveals that oscillations persist up to 350 mK and enables us to estimate the phase coherence length of holes to be $L_\varphi = 2 \, \mu$m at the base temperature of 65 mK. Although this value is one order of magnitude smaller than the typical values for $L_\varphi$ in electron systems, it shows that fabrication of phase-coherent p-type GaAs nanodevices is within the reach of present nanofabrication technologies.
Appendices

A List of samples

<table>
<thead>
<tr>
<th>name</th>
<th>picture</th>
<th>wafer</th>
<th>structure</th>
<th>measurements</th>
</tr>
</thead>
<tbody>
<tr>
<td>QPC</td>
<td><img src="image" alt="QPC picture" /></td>
<td>Bochum12029</td>
<td>165nm wide Quantum Point Contact</td>
<td>chapter 7</td>
</tr>
<tr>
<td></td>
<td></td>
<td>45nm deep 2DHG</td>
<td>$N = 3.8 \times 10^{11} \text{ cm}^{-2}$</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>$\mu = 200'000 \text{ cm}^{2}/\text{Vs}$</td>
<td></td>
</tr>
<tr>
<td>Quantum dot</td>
<td><img src="image" alt="Quantum dot picture" /></td>
<td>Bochum12029</td>
<td>Rectangular quantum dot with lithographic dimensions</td>
<td>chapter 8</td>
</tr>
<tr>
<td></td>
<td></td>
<td>45nm deep 2DHG</td>
<td>$N = 3.8 \times 10^{11} \text{ cm}^{-2}$</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>$\mu = 200'000 \text{ cm}^{2}/\text{Vs}$</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>$430 \times 170 \text{ nm}^{2}$</td>
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</tr>
<tr>
<td>Large ring</td>
<td><img src="image" alt="Large ring picture" /></td>
<td>Bochum12029</td>
<td>Quantum ring with an orbital radius of 420nm</td>
<td>chapter 9</td>
</tr>
<tr>
<td></td>
<td></td>
<td>45nm deep 2DHG</td>
<td>$N = 3.8 \times 10^{11} \text{ cm}^{-2}$</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>$\mu = 200'000 \text{ cm}^{2}/\text{Vs}$</td>
<td></td>
</tr>
<tr>
<td>Small ring</td>
<td><img src="image" alt="Small ring picture" /></td>
<td>Bochum12029</td>
<td>Quantum ring with a radius of the mean orbit of 160nm and lithographic radius of 320nm</td>
<td>chapters 8, 9</td>
</tr>
<tr>
<td></td>
<td></td>
<td>45nm deep 2DHG</td>
<td>$N = 3.8 \times 10^{11} \text{ cm}^{-2}$</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>$\mu = 200'000 \text{ cm}^{2}/\text{Vs}$</td>
<td></td>
</tr>
</tbody>
</table>
### B Processing of GaAs samples in FIRST

Old numbers without the 40% filter on the MJF3 are given in parentheses.

<table>
<thead>
<tr>
<th>Process</th>
<th>Description of process steps</th>
<th>Remarks</th>
</tr>
</thead>
</table>
| Cleaning sample (first time only!) | Acetone 1 min ultrasound max. power  
Isopropanol 1 min ultrasound (tilt sample when removing from recipient)  
H₂O  
Dry with N₂  
Heat 2 min at 115°C  
HCl-dip for 8 sec | Christoph: sometimes 2-3 min may use H₂O with ultrasound  
Roland: 90°C is enough only for dirty samples |
| Lithography for mesa (positive process) | Store a few 10th of a ml of Shipley S1805 photoresist in a disposable pipette  
Put 1 drop on sample and immediately spin on: 5000rpm for 60 sec  
Soft bake: hotplate at 115°C for 2 min  
Exposure for edge removal: 20(5...10) sec soft contact  
Develop: Shipley MF319 for 20 sec  
Rinse thoroughly with H₂O for at least 2 min, dry with N₂  
Exposure: 4(1.8) sec hard contact  
Develop: Shipley MF319 for 22...30(20...23) sec | wait till pipette content is at room temperature  
if faster: problems with sticking, but shallower resist edges  
Christoph: let cool 1 min after bake  
if necessary, wash mask prior to use  
Davy: Additional edge removal Exp. 10 sec Dev. 30 sec  
if necessary, wash mask prior to use  
control optically |
| Lithography for backgate (using the optical microscope) | Hoechst AZ5214E  
Spin on:2000 rpm for 3 sec, 4500...5000 rpm for 45 sec  
Soft bake: hotplate 90°C for 2 min  
Use the aperture to reduce the spot size, set the intensity of the bulb to the maximum, remove the UV filter and all polarizers to start the exposure, reinsert UV filter to stop, reduce the intensity to barely visible | wait till pipette content is at room temperature  
put drop during first 3 sec  
for the 100x magnification |
| Mesa/Backgate etch | When thinning down acids, always pour in water first: Remember: “A” comes before “W” in the alphabet.  
"Nie das Wasser in die Säure, sonst passiert das Ungeheure.”  
"Assis dans l’eau.”  
If necessary, use stirring staff/magnetic stirrer  
Note: delayed etching: first 30 sec smaller etch rate, second 30 sec larger etching rate  
H₂O+H₂SO₄+H₂O₂ = 100 : 3 : 1 (50 ml:1.5 ml:0.5 ml),  
Etch rate: 45 sec≈45 nm(Stefan:40 sec≈60 nm at T = 22.7°C)  
Mesa : typically 80 nm ± 2 min  
Backgate : typically 480...1000 nm ± 8...10 min | Christoph: use 100:3:3 for Backgate (6.5 min±800 nm)  
Stop in water, dry with N₂ |
| Cleaning sample (after every resist depositing step) | Acetone 1 min ultrasound 4/9 power  
Isopropanol 1 min ultrasound (tilt sample when removing from recipient)  
H₂O  
Dry with N₂ | Christoph: “first time cleaning” also possible |
| Lithography with Hoechst AZ5214E (for ohmic contacts, negative process) | Spin on: 2000 rpm for 3 sec, 5000 rpm for 45 sec, photoresist: Hoechst AZ5214E  
Soft bake: hotplate at 90°C for 2 min  
Exposure: hard contact 2(1) sec  
Hard bake: hotplate at exactly 115°C for 2 min  
Exposure of everything (without mask): soft contact 20(10) sec  
Develop: Shipley MF319 for 90 sec  
Rinse thoroughly with H₂O for at least 2 min, dry with N₂ | let cool down 1 min  
control optically |
<table>
<thead>
<tr>
<th>Process</th>
<th>Description of process steps</th>
<th>Remarks</th>
</tr>
</thead>
<tbody>
<tr>
<td>Cleaning before Deposition</td>
<td>Oxygen plasma ash for 1...3 min HCl dip etch 5 sec, stop in water</td>
<td>note: etching for 60 min reduces mobility by a factor of 4 can avoid bubbles on ohmic contacts</td>
</tr>
<tr>
<td>Contact metallization</td>
<td>Ask Christoph OR Urszula OR Hanauer Scherrer (Tel. 01/63-32072 or -33261 or easiest by e-mail: <a href="mailto:scherrer@phys.ethz.ch">scherrer@phys.ethz.ch</a>). Use process &quot;2nd ohmic contact&quot; (Ge/Au/Ge/Au/Ni/Au = 18/50/18/50/40/100 mm)</td>
<td>important: evaporate contacts soon after plasma and/or dip etch</td>
</tr>
<tr>
<td>Lift-off</td>
<td>Immerse sample in acetone and wait for 5 min...4 h (let the acetone creep under the photosresist) If necessary, use ultrasound at lowest power (1/9) for a few seconds, then wait again Take sample out carefully. Avoid gold particles sticking on the surface by e.g. rinsing with acetone while removing the sample</td>
<td>for critical lift-off use boiling Acetone and watergun for removing</td>
</tr>
<tr>
<td>Ohmic Contact anneal</td>
<td>Let it dry &amp; outgas at 120°C for 1 min Annealing with inert gas: 5...40 sec at 450°C, inert gas -0.5 bar</td>
<td>Christoph: 430°C / 30s, probably shorter time is better</td>
</tr>
<tr>
<td>BG Contact anneal</td>
<td>Annealing with inert gas: 30 sec at 450°C (not more!), inert gas -0.5 bar</td>
<td>somehow better to anneal BG and 2DEG at the same time</td>
</tr>
<tr>
<td>Bondpad reinforcement</td>
<td>Perform lithography with Hoechst AZ5214E (for ohmic contacts, negative process) using the bondpad reinforcement mask Oxygen plasma ash for 1 min. Contact metallization: Use 100 nm Au.</td>
<td>Align properly, otherwise contacts lost due to etching!</td>
</tr>
<tr>
<td>Positive process lithography for topgate fingers (not recommended)</td>
<td>Spin on: 2000 rpm for 3 sec, 5000 rpm for 45 sec, photosresist: Hoechst AZ5214E Soft bake: hotplate at 90°C for 2 min Exposure: negative mask with topgate fingers in hard contact for 4...6(2) sec Develop: Shipley MF319 for 90 sec Rinse thoroughly with H₂O₂ for at least 2 min, dry with N₂</td>
<td>Shipley with 4.5(1.8) sec / 25 sec gives finer fingers control optically</td>
</tr>
<tr>
<td>Negative process lithography for topgate fingers</td>
<td>Spin on: 2000 rpm for 3 sec, 5000 rpm for 45 sec, photosresist: Hoechst AZ5214E Soft bake: hotplate at 90°C for 2 min Exposure: hard contact 3.8.5(0.8) sec Hard bake: hotplate at exactly 115°C for 2 min Exposure of everything (without mask): soft contact 20...25(10) sec Develop: Shipley MF319 for 70...90(90) sec Rinse thoroughly with H₂O₂ for at least 2 min, dry with N₂</td>
<td>Shipley with 4.5(1.8) sec / 25 sec gives finer fingers control optically</td>
</tr>
<tr>
<td>Top gate finger metallization</td>
<td>Do as in contact metallization, using Ti/Au=10/80 nm</td>
<td></td>
</tr>
<tr>
<td>Lithography for e-beam</td>
<td>Spin on: PMMA 4:5 with 6000 rpm for 60 sec (→ 90 nm)</td>
<td>Christoph: 2 layers for bigger things: P(MMA/MAA); PMMA.</td>
</tr>
<tr>
<td>Etching of AFM-defined samples using HCl</td>
<td>Create sample using AFM lithography. Consider increase in depletion length (by approximately 10 nm in backgated samples). Etch for 10 sec in HCl (conc.), stop with H₂O.</td>
<td></td>
</tr>
<tr>
<td>General remarks</td>
<td>When switching off the mask aligner, let the lamp cool down for 30 min before turning off N₂ If chemicals are missing: ask FIRST team. Empty bottles have to be purged with water and placed in the lock (B125.1).</td>
<td></td>
</tr>
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</table>
Publications

Magnetotransport in C-doped AlGaAs heterostructures
Boris Grbić, Christoph Ellenberger, Thomas Ihn, Klaus Ensslin, Dirk Reuter, and Andreas D. Wieck

Single-hole transistor in p-type GaAs/AlGaAs heterostructures
Boris Grbić, Renaud Leturcq, Klaus Ensslin, Dirk Reuter, and Andreas D. Wieck

Electronic properties of C-doped (100) AlGaAs heterostructures
Boris Grbić, Christoph Ellenberger, Thomas Ihn, Klaus Ensslin, Dirk Reuter, and Andreas D. Wieck

Hole transport in p-type GaAs quantum dots and point contacts
Boris Grbić, Renaud Leturcq, Thomas Ihn, Klaus Ensslin, Dirk Reuter, and Andreas D. Wieck

Aharonov-Bohm oscillations in the presence of strong spin-orbit interaction
Boris Grbić, Renaud Leturcq, Thomas Ihn, Klaus Ensslin, Dirk Reuter, and Andreas D. Wieck
submitted, cond-mat:0704.1264

Aharonov-Bohm oscillations in p-type GaAs quantum rings
Boris Grbić, Renaud Leturcq, Thomas Ihn, Klaus Ensslin, Dirk Reuter, and Andreas D. Wieck
submitted to Physica E

Strong spin-orbit interaction in carbon doped p-type GaAs heterostructures
B. Grbić, R. Leturcq, T. Ihn, K. Ensslin, D. Reuter, and A. D. Wieck
submitted to Physica E
Bibliography


[27] T. Ihn, Halbleiter-Nanostrukturen (Lecture notes ETH Zurich, 2005).


[51] The 1D Poisson-Schrödinger simulation is taken from the G. Snider’s web-site http://www.nd.edu/ gsnider/1DPoisson.hqx (????).


Acknowledgements

Working on my PhD thesis at ETH Zurich was indeed unique experience - this is partly due to the great infrastructure ETH provides for conducting top-level research, but primarily due to stimulative and truly international scientific environment at ETH. Now, in the end of my PhD way I would like to express my sincere gratitude to those people who contributed to my PhD work and my development as a physicist, and who made my time in Zurich unforgettable.

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I learned a lot from Barbara Simović about how to approach physics experiments. Discussions with her and her comments often steered me in the right direction - thanks a lot for that. She was also the master to find some intriguing topics to chat about during our lunch breaks.

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Paul Studerus is a true magician in electronics - one needs only to imagine certain filter or amplifier and the desired "box" will soon be ready on the table. Beside the fact that our experiments would not have been possible without Paul's huge expertise in electronics, I want also to acknowledge his readiness to make electronics closer to us physicists. Cecil Barengo deserves my thanks for making all kinds of mechanical gadgets we need in our experiments.

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Although his last name does not start with G, Lorenz Meier was like a member of the G4 gang. Thanks to Lorenz for showing us in practice how does the Swiss democracy work by initiating on-line group votings about all imaginable group issues. Also the visit of the nuclear power plant which he organized was indeed a unique experience.

I am grateful to Thomas Müller for giving me the opportunity to boost my sport-ego by beating him in a couple of table-tennis sets during the breaks in the lab. It was always fun to exchange teasing comments with him.

I acknowledge Andreas Pfund for his efforts to master the pronunciation of my last name. He has almost managed to pronounce the word Grbić without adding some
more vowel-like sounds to it - amazing!

The B19 crew, Slavo Kičin, Magdalena Hüfner and Alessandro Pioda, is gratefully acknowledged for being very generous whenever I came to their well ordered lab to take some piece of equipment for my measurements.

I wish best of luck to my successors on the p-type GaAs project, Yashar Komijani and Miklós Csontos. Keep being excited with holes!

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## Curriculum Vitae

<table>
<thead>
<tr>
<th>Date</th>
<th>Event</th>
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<tbody>
<tr>
<td>09/1984-06/1992</td>
<td>Primary School &quot;Vuk Karadžić&quot; in Belgrade, Serbia</td>
</tr>
<tr>
<td>09/1992-06/1996</td>
<td>Mathematical High School in Belgrade, Serbia</td>
</tr>
<tr>
<td></td>
<td>Specialized High School for students gifted in mathematics and</td>
</tr>
<tr>
<td></td>
<td>natural sciences - only school of its kind in Serbia</td>
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<tr>
<td>1995, 1996</td>
<td>Honourable mentions awarded at</td>
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<tr>
<td></td>
<td>26th International Physics Olympiad in Canberra, Australia</td>
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<tr>
<td></td>
<td>27th International Physics Olympiad in Oslo, Norway</td>
</tr>
<tr>
<td>10/1996-02/2002</td>
<td>Faculty of Physics - University of Belgrade, Serbia</td>
</tr>
<tr>
<td></td>
<td>Diploma thesis: &quot;Bands of quasiparticle states in superconducting</td>
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<tr>
<td></td>
<td>superlattices&quot;; Supervisor Prof. Zoran Radović</td>
</tr>
<tr>
<td>10/2000-12/2000</td>
<td>International student exchange - National Institute</td>
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<tr>
<td></td>
<td>of Industrial Technology (INTI), Buenos Aires, Argentina</td>
</tr>
<tr>
<td>10/1996-06/2002</td>
<td>Teaching associate for physics in Mathematical High School</td>
</tr>
<tr>
<td></td>
<td>in Belgrade, Serbia</td>
</tr>
<tr>
<td>02/2002</td>
<td>Award from Serbian Royal Family for the top 100 students in Serbia</td>
</tr>
<tr>
<td>September 2002</td>
<td>Start of PhD research work in the Solid State Physics Laboratory</td>
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<tr>
<td></td>
<td>(ETH Zurich) in the Nanophysics group of Prof. Klaus Ensslin</td>
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