Stellingen

behorende bij het proefschrift

*Electron Transport in Mesoscopic GaAs/AlGaAs structures
with Superconducting Contacts*

door Kars-Michiel Lenssen

I

Met gediffundeerde contacten kunnen supergeleidende effecten in een tweedimensionaal elektronengas geïnduceerd worden.

Dit proefschrift, hoofdstuk 7, 8 en 9

II

Het is mogelijk systemen van tweedimensionaal gekoppelde quantum dots te realiseren met behulp van gatestructuren.

Dit proefschrift, hoofdstuk 3 en 4

III

Het verdient aanbeveling de achtergronden in spectra verkregen met behulp van low-energy ion scattering (LEIS), meer op de voorgrond te laten treden.

K.-M.H. Lenssen, afstudeerverslag T.U. Eindhoven (1989)

IV

Het tweedimensionale elektronengas (2DEG) op vloeibaar helium is een zeer geschikt modelsysteem voor de bestudering van o.a. het Hall-effect in een niet-gedegeneriseerd 2DEG.

V

Met behulp van temporele decompositie kan spraak zeer zuinig gecodeerd worden, terwijl de verstaanbaarheid behouden blijft.

K.-M.H. Lenssen, IPO-rapport 629 (1988)

VI

In netwerken van Josephson-juncties kan een analogon van Andreev-reflectie bestaan voor vortices.

VII

Het diagram van de bandbuiging in een metaal/ n^+ -GaAs/ n -GaAs-systeem dat Braslau gebruikt is onjuist, hoewel het later vele malen gekopieerd is.

N. Braslau, J. Vac. Sci. Technol. **19**, 803 (1981),

N. Braslau, Thin Solid Films **104**, 391 (1983),

M. Shur, GaAs Devices and Circuits, Plenum Press, New York (1987)

VIII

Fundamenteel mesoscopisch onderzoek kan commercieel interessant zijn door sample leasing in te voeren.

IX

Het verdient geen aanbeveling de term nanolithografie in het Nederlands te vertalen, aangezien dit verkeerde associaties zou kunnen oproepen.

X

Een ervaren blikjesverzamelaar in Central Park (New York) verdient op sommige dagen meer dan een onderzoeker in opleiding.

XI

Het feit dat de term miniaturisering uit de kloosterschilderkunst afkomstig is, is te herkennen doordat het nog steeds monnikenwerk is.

XII

De reclamespotjes voor Jif-allesreiniger (met microkracht) en Nivea-crème (met nanosferen) doen de vraag rijzen of nanostructuren vercommercialiseren.

Delft, 9 mei 1994

**Electron Transport in Mesoscopic
GaAs/AlGaAs Structures
with Superconducting Contacts**

Electron Transport in Mesoscopic GaAs/AlGaAs Structures with Superconducting Contacts

PROEFSCHRIFT

ter verkrijging van de graad van doctor
aan de Technische Universiteit Delft,
op gezag van de Rector Magnificus
Prof. ir. K. F. Wakker,
in het openbaar te verdedigen
ten overstaan van een commissie,
door het College van Dekanen aangewezen, op

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Kars-Michiel Hubert LENSSEN

natuurkundig ingenieur,
geboren te Utrecht.

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De hele mesoscopische fysica is een "zero-bias anomaly".

(C. W. J. Beenakker)

*A physical understanding is a completely unmathematical, imprecise,
and inexact thing, but absolutely necessary for a physicist.*

(R. P. Feynman, *The Feynman lectures on physics vol. 2*)

Be patient, for the world is broad and wide.

(E. A. Abbott, *Flatland*)

Preface

This thesis reports the results of a graduate research, carried out in the group Solid State Physics/Superconductivity of Prof. Dr. Ir. J. E. Mooij at Delft University of Technology in the Netherlands. The aim was to combine semiconductor nanostructures with superconducting contacts in order to study new mesoscopic transport phenomena in such systems. Such phenomena have indeed experimentally been observed, as will be described here. This will be preceded by a theoretical introduction, a numerical theoretical study and a description of the realization of the new type of superconducting contacts. Besides that, measurements of electron transport in a two-dimensional lateral superlattice will be discussed in the first part of the thesis.

Several parts of this thesis have been or will be published elsewhere in a slightly different form:

- chapter 6: Phys. Rev. B **44**, 470 (1991)
- chapter 8: Appl. Phys. Lett. **63**, 2079 (1993)
- chapter 9: *Editions Frontières*, Proc. on “Coulomb and interference effects in small electronic structures”, to be published
- chapters 3, 4 and 10: prepared for publication.

At this place I would like to thank Prof. Dr. Ir. J. E. Mooij and my advisor Dr. C. J. P. M. Harmans for giving me the opportunity to do this research in a group at the frontier of science and for their scientific and mental support. The undergraduate students M. Matters, M. E. J. Boonman, L. A. Westerling and P. C. A. Jeekel have contributed significantly to the scientific results and the pleasant working atmosphere. I am grateful to A. van der Enden, L. Lander and P. Nouwens for technical assistance.

Because of the multidisciplinary character this research wouldn't have been successful without the great cooperation with other scientists. In the first place I would like to thank Prof. Dr. J. H. Wolter and his group at Eindhoven University of Technology, in particular Dr. M. R.

Leys, for providing us with high-mobility 2DEG's and for a very pleasant and fruitful cooperation. Also I am very grateful to Prof. Dr. C. W. J. Beenakker of the University of Leiden for many valuable discussions and theoretical support. The people and the facilities of the Delft Institute of Micro-Electronics and Submicron technology, in particular in the group of Prof. Dr. Ir. S. Radelaar, were essential for the realization of the nanostructures. I also thank the Philips Research Laboratories, and Dr. M. Holland and Prof. Dr. S. P. Beaumont of the University of Glasgow for providing us with MBE-grown heterostructures.

I have benefited from many interactions with great scientists. In this respect I would like to thank Dr. J. P. Eisenstein, Dr. A. Kastalsky, Dr. A. W. Kleinsasser and Prof. Dr. H. Lüth for their hospitality. Discussions with them, Prof. Dr. H. van Houten, Dr. Z. G. Ivanov, Prof. Dr. Ir. T. M. Klapwijk and his group (in particular Dr. J. R. Gao and Dr. Ir. B. J. van Wees), and Dr. H. Takayanagi are gratefully acknowledged. I thank Prof. Dr. I. O. Kulik for sending me a rare copy of a Russian preprint. Finally, I gratefully acknowledge the financial support of the Organization for Fundamental Research on Matter (FOM), which was a very pleasant and good employer.

Kars-Michiel Lenssen
Delft, March 1994

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Chapter 1

General introduction

1.1 Introduction

Superconductivity has been discovered in 1911, semiconductor physics exists about half a century, mesoscopic physics a little more than a decade. This thesis describes one of the first investigations meant to combine all these science fields, a study of mesoscopic semiconductor structures with superconducting contacts.

The first part deals with an example of semiconductor nanostructures. After that, the coupling of superconductors and semiconductors is described. The last part discusses measurements in mesoscopic superconductor/semiconductor systems.

1.2 Mesoscopic physics in semiconductors

Mesoscopic physics studies systems of which the size in at least one dimension is of the same order as a physically relevant length. This is a rather general, but in practice commonly used definition. For our purpose we only consider length scales that characterize electronic transport.

Dependent on which physical length is chosen as a reference, different transport regimes can be distinguished. This is schematically illustrated in figure 1.1. Large systems can be described without the use of quantum mechanics; this is the classical regime of the macroscopic world. If however the size L becomes smaller than the phase-coherence length l_φ , the wave properties of electrons can not be ignored. Over a distance l_φ phase information of the electrons is conserved, which causes inter-

Figure 1.1: Overview of typical length scales associated with different regimes of electronic transport (the symbols are explained in the text)

ference effects. Here the field of mesoscopic transport starts with the diffusive regime. For distances below the elastic mean free path l_e electrons are not scattered. This means that their paths can be deduced from the initial position and velocity: the electrons move ballistically. In this regime electrons can often be described as billiard balls. At even smaller sizes the wavelength λ of the electrons becomes important. Now the electrons have to be treated fully quantum mechanically; in this regime quantization effects (e.g. of conductance) occur, sometimes very reminiscent of phenomena in wave mechanics and optics.

At further decreasing L many theoretic concepts of solid state physics (like band structure) lose their validity, if the atomic distance d is approached. However, here we are leaving the field of mesoscopics and entering the microscopic world.

Other typical lengths can also be important. Magnetic fields can introduce length scales, which can replace the relevant length without applied field (e.g. the magnetic length l_m). By inducing a periodic potential in a system also this period becomes a relevant unit.

As the reader may have noted, mesoscopic systems often allow the use of textbook physics. This is one of the charms of mesoscopic physics: often it gives the possibility to study model systems (like a particle in a box, modes in a wave guide, ballistic motion, or even the formation of an artificial band structure) [1].

Since mesoscopic effects are best observed if the sizes in all dimensions are “small”, low-dimensional systems are preferred. A well-known

example, which will be the basic system in this thesis, is the two-dimensional electron gas (2DEG) which can be formed at an interface between GaAs and AlGaAs. The band diagram of such a semiconductor heterostructure is shown in figure 1.2. Since AlGaAs has a larger band

Figure 1.2: *Schematic band diagram of a GaAs/AlGaAs heterostructure with a two-dimensional electron gas*

gap than GaAs a quantum well is formed at the interface, if the AlGaAs is *n*-doped. Electrons in this quantum well can freely move in the plane of the interface, but perpendicular to the interface two-dimensional subbands are formed. If the doping is chosen such that only the lowest subband is occupied by electrons, a (quasi-)two-dimensional electron gas is obtained. To achieve a high mobility 2DEG the interface has to be nearly perfect. By means of advanced MBE-techniques¹ this can be realized, since GaAs and AlGaAs have almost the same lattice spacing.

Using gates electrons can be confined also in the directions of the interface (see figure 1.3). If a negative voltage is applied to a metal film on top of the heterostructure (the gate), a tuneable potential barrier is induced in the underlying 2DEG. Large negative gate voltages can completely deplete 2DEG regions.

By patterning the gates many interesting structures can be fabricated: constrictions, cavities, grids, etc. Electron beam lithography makes it possible to define structures with sizes comparable to the electron wavelength in a thin resist layer on top of the heterostructure. Parts of this resist that are exposed to the electron beam change their

¹MBE is Molecular Beam Epitaxy

Figure 1.3: *Schematic view of the action of a gate (cross-sectional view); the dashed line shows the boundary of the induced potential barrier*

composition and are removed during development. The realized pattern can be transferred to a metal film by a lift-off technique. First, a metal film (e.g. Ti/Au) is evaporated on top of the patterned resist layer. When the resist is now dissolved (in warm acetone), all metal will come off except in those areas where the metal is in direct contact with the semiconductor.

1.3 Superconductor/semiconductor systems

Experimentally it is known that current can flow without resistance (i.e. a supercurrent) from one superconductor via a normal material to another superconductor, if the region of normal material is thin enough. This system is called a weak link [2]. There has always been a large interest in combining superconductors with semiconductors, since this could provide the possibility to modulate superconducting properties by varying the carrier density in the semiconductor. In a weak link consisting of two superconductors coupled by a semiconductor the supercurrent may be controlled by a gate voltage. Such a three-terminal device could act as a superconducting transistor [3].

Besides the potential device applications superconductor/semiconductor systems are also interesting from a fundamental point of view. Especially when semiconductor nanostructures could be combined with superconducting electrodes many new physical phenomena can be expected. For example, it has recently theoretically been predicted that the critical supercurrent through a quantum point contact will be quantized [4, 5]. But also interference effects in mesoscopic systems will be drastically influenced by the presence of superconductors, because reflection from a clean superconducting boundary is very different from

normal reflection [6]. This expected field of challenging new physics was the motivation for our research.

Since the 2DEG in a GaAs/AlGaAs heterostructure has proven to be a very suitable system for the study of normal mesoscopic effects, we also wanted to use it in combination with superconducting contacts. The realization of such contacts is however far from trivial, because the 2DEG is situated typically $0.1\text{ }\mu\text{m}$ below the surface of the heterostructure. For this purpose a new process for the fabrication of diffused superconducting contacts had to be developed, which could be combined with gate structures. The observed new physical phenomena, however, make this effort certainly worth while.

1.4 Scope of the thesis

In this thesis first a few theoretical concepts in normal mesoscopic physics will be introduced (chapter 2). After that, measurements on a sample with a grid-shaped gate, realized using state-of-the-art electron beam lithography, will be treated. Electronic transport properties of this two-dimensional superlattice will be discussed in weak (chapter 3) and strong magnetic fields (chapter 4). Signatures of phase-coherent coupling of respectively classical and quantum mechanical electron orbits (giving rise to band-structure effects) will be presented.

Then semiconductor nanostructures will be combined with superconductors. A short description of models of interfaces between superconductors and normal materials (chapter 5) is followed by a numerical application of a developed generalized transmission formalism for supercurrent flow in superconductor/semiconductor systems (chapter 6). Next the realization and characterization of a new type of diffused superconducting contacts to the 2DEG in GaAs/AlGaAs heterostructures is described (chapter 7). In chapter 8 the observation of Andreev reflection at these contacts and the influence of gates on this superconducting effect are shown. The last two chapters discuss new phase-coherent effects induced by superconducting contacts, e.g. in combination with a quantum point contact: conductance quantization in enhanced units, reflectionless tunneling and enhanced weak localization (chapters 9 and 10). The thesis is concluded with summaries in English and Dutch.

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Chapter 2

Theoretical concepts in mesoscopic physics

2.1 Introduction

In the field of mesoscopic physics in semiconductors many good publications exist (see, for example Refs. [1–3] and references therein). Therefore we will only discuss a few relevant concepts that will be encountered later in this thesis. The first one is an important interference effect: weak localization. After this phenomenon which is based on elastic scattering, we will turn to the ballistic regime with the conductance quantization in a quantum point contact. This is an example of confinement in two directions; however, electrons can also be confined in all three dimensions: in quantum dots. Of course, it is challenging to combine several of these artificial “atoms” to arrays of coupled quantum dots, in order to study artificial “crystals”. Since this intriguing research is the subject of chapter 3 and 4, artificial band structures will be introduced in the last section of this chapter.

2.2 Weak localization

Since purely elastic scattering does not destroy phase coherence (in general the phase-coherence length l_φ is larger than the elastic mean free path l_e), quantum interference effects can occur in ballistic, but also in diffusive transport. In the Feynman path description [4] the probability $P(\vec{r}, \vec{r}', t)$ for motion from point $\vec{r}$ to point $\vec{r}'$ in a time t can be expressed in the probability amplitudes A_i , corresponding with all

trajectories from $\vec{r}$ to $\vec{r}'$ of duration t (see figure 2.1a):

Figure 2.1: Schematic picture of the Feynman path description for a) general trajectories and b) backscattering trajectories

$$P(\vec{r}, \vec{r}', t) = \left| \sum_i A_i \right|^2 = \underbrace{\sum_i |A_i|^2}_{\text{classical}} + \underbrace{\sum_{i \neq j} A_i A_j^*}_{\text{interference}}. \quad (2.1)$$

The right-hand side consists of a term equal to the classical diffusion probability plus a term due to quantum interference.

In general the phases of the numerous, different trajectories are uncorrelated and thus the interference term averages out. However, for the specific case where $\vec{r} = \vec{r}'$ (backscattered trajectories, see figure 2.1b), the probabilities A^+ and A^- for clockwise and counterclockwise propagation around the closed path are identical because of time-reversal invariance:

$$A^+ = A^- \equiv A \quad (2.2)$$

and so

$$P(\vec{r}, \vec{r}, t) = 2A^2 + 2A^2. \quad (2.3)$$

The coherent backscattering probability is twice the classical result. This is the essence of weak localization: enhanced probability of backscattering [5]. It will be clear, that weak localization is suppressed by the application of a weak magnetic field, since in that case time-reversal invariance is broken.

2.3 Conductance quantization

By means of a split-gate structure (see figure 2.2 and also figure 1.3) on top of a GaAs/AlGaAs heterostructure a short and narrow constriction can be made in a two-dimensional electron gas (2DEG). The width W

Figure 2.2: Schematic top view of a quantum point contact; W is typically 250 nm, the hatched areas are gates, the dashed lines represent the boundary of the induced electrostatic barrier

of the constriction is of the order of the Fermi wavelength ($\lambda_F \approx 40$ nm) and can be varied by modulation of the negative voltage on the gates. Measurements of the conductance of such a system as a function of W have led to the discovery of conductance quantization [6, 7].

Since the dimensions of a quantum point contact are much smaller than the elastic mean free path ($l_e \approx 10$ μm) the electron transport is ballistic. Classically one would expect that the conductance G is proportional to W :

$$G_{\text{class}} = \frac{2e^2}{h} \cdot \frac{2W}{\lambda_F}. \quad (2.4)$$

However, quantum effects cannot be ignored. The transverse momentum $\hbar k_y$ only can have discrete values $\pm n\pi/W$ (with $n = 1, 2, 3, \dots$). Each time the width is increased by an amount $\lambda_F/2$ an extra mode or 1D subband becomes available for transport. The total number of subbands below the Fermi level is given by the number of half the electron wavelength that fits within W , called N .

The amount of current that a subband can carry, in case of a small difference $\delta\mu$ between the chemical potential at both sides of the quantum point contact, is given by:

$$I_n = ev_n \rho_n \delta\mu, \quad (2.5)$$

with v_n the group velocity and ρ_n the density of states (with a group velocity of the relevant direction). Substitution in this expression of $v_n = \frac{\partial \varepsilon_n}{\hbar \partial k_x}$ and $\rho_n = \left(\pi \frac{\partial \varepsilon_n}{\partial k_x} \right)^{-1}$, assuming a dispersion relation $\varepsilon_n(k_x)$ corresponding with the n th subband, yields

$$I_n = \frac{2e}{h} \delta\mu, \quad (2.6)$$

independent of the subband index n , since all mode dependent factors cancel! Using the definition of the conductance

$$G \equiv \frac{Ie}{\delta\mu} = \frac{2e^2}{h} \cdot N, \quad (2.7)$$

it can now easily be seen that the conductance is quantized in units of $2e^2/h$.

This semiclassical result is confirmed by a quantummechanical treatment based on the Landauer formula [8]

$$G = \frac{2e^2}{h} \text{Tr } \mathbf{t} \mathbf{t}^\dagger, \quad (2.8)$$

with $\mathbf{t}$ the $n \times n$ transmission matrix. For a ballistic quantum point contact $\mathbf{t}$ is the unity matrix, so that equation (2.7) is obtained again.

2.4 Artificial band structure

In a crystal the periodic potential of the lattice causes energy gaps to appear in the band structure. Already in 1970 it was proposed [9] that artificial energy gaps could be created by superimposing an artificial periodic modulation of the crystal. It was suggested that such a superlattice could be obtained by growing epitaxially alternating layers of different semiconductors. Since the period a of this superlattice is (much) larger than the crystal lattice spacing d , this will induce many narrow bands (minibands) within the conduction band, separated by small energy gaps (minigaps). The band structure of such a superlattice is shown in figure 2.3.

Since the superlattice formed by alternating layers is one-dimensional, the density of states is not zero in the minigaps because of the free motion in the plane of the layers. Lateral superlattices, formed by a periodic gate structure on top of a 2DEG system, provide the possibility to vary

Figure 2.3: *Schematic band structure of a superlattice (solid lines); the dashed line represents the original conduction band of the crystal*

the potential in two directions, thus achieving true minigaps of zero density of states. A two-dimensional periodic potential can be induced by a gridshaped gate, as will be studied in chapters 3 and 4.

For a sinusoidal potential the first minigap is centred around $E_1 \approx \frac{\hbar^2 \pi^2}{2ma^2}$. It is lithographically not yet achievable to make the period a so small that E_1 becomes of the order of the Fermi energy. Thus in this way the band structure can't be examined. There is, however, another way to study the dynamics of lateral superlattices: by applying a perpendicular magnetic field B . In that case there are two relevant length scales: the period a of the superlattice potential and the magnetic length $l_m \equiv \sqrt{\hbar/eB}$. We can distinguish two limits.

In one limit, the Landau regime, the periodic potential is regarded as a weak perturbation on the usual magnetic Landau-level structure. If the magnetic flux through one unit cell of the periodic potential (with area S) is given by $\phi = eSB/\hbar = p/q$ (with p and q integers), each Landau level splits into p subbands of equal degeneracy. This will result in oscillations periodic in $1/B$ in the magnetoconductance. These so-called Weiss oscillations, which can be considered in terms of commensurability of the cyclotron radius with the superlattice period, are recently observed in weak periodic potentials [10]. Measurements on

a two-dimensional lateral superlattice in this regime are discussed in chapter 3.

In the opposite limit, the Onsager regime, the magnetic field is treated as a perturbation on the periodic potential. In this case the magnetotransconductance is expected to be periodic in B . If the magnetic length l_m is of the same order as a an exotic energy spectrum is theoretically predicted for a two-dimensional potential [11]. In natural lattices the required magnetic fields are gigantic, but in lateral superlattices this situation can be obtained. The observation of these effects requires phase coherence of the electrons over several unit cells. In a one-dimensional lateral superlattice (consisting of 15 cavities in series) the experimental observation of a magnetically induced band structure has already been claimed [12]. Measurements on a two-dimensional lateral superlattice in high magnetic fields are described in chapter 4 of this thesis.

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Chapter 3

Observation of quantum interference in a phase-coherent two-dimensional superlattice

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Abstract

Electron transport in a two-dimensionally periodically modulated potential has been studied by means of gate structures on top of a GaAs/AlGaAs heterostructure. In low magnetic fields ($B < 1$ T) conductance minima are measured due to localization of electron orbits. Superimposed on these minima Aharonov-Bohm oscillations are observed with periods in B which are submultiples of that corresponding to a single localized orbit. This new phenomenon is attributed to phase-coherent coupling of localized orbits.

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3.1 Introduction

Artificial lateral superlattices provide a convenient model system for the research of a number of fundamental effects in electron transport. Employing a grid-shaped metallic surface gate to define the superlattice and by tuning the voltage on this gate, different transport regimes can be obtained: from a weak periodical potential via an artificial two-dimensional crystal to an array of weakly coupled quantum dots. Each of these regimes has typical physical phenomena (theoretically predicted or already measured), ranging from the so-called Weiss oscillations [1], via band structure effects [2] to Coulomb effects. In particular effects due to two-dimensional coupling of quantum dots are still barely experimentally studied. Here we will discuss measurements in low magnetic fields [3] in the regime of a weak potential, i.e. in a two-dimensional array of strongly coupled quantum dots.

3.2 Sample fabrication

For the samples the two-dimensional electron gas (2DEG) in a GaAs/AlGaAs heterostructure is used, because of its high mobility, the large Fermi wavelength and its excellent gating properties. The 2DEG has a mobility $\mu_e \gtrsim 100 \text{ m}^2/\text{Vs}$ and an electron density $n_s \approx 1.3 \cdot 10^{15} \text{ m}^{-2}$, resulting in a Fermi wavelength $\lambda_F \gtrsim 70 \text{ nm}$ and an elastic mean free path $l_e \gtrsim 6 \text{ }\mu\text{m}$. The material was grown by Molecular Beam Epitaxy and consists of a GaAs substrate with a GaAs/AlGaAs superlattice, a $0.5 \text{ }\mu\text{m}$ undoped GaAs layer, a 41 nm undoped $\text{Al}_{0.33}\text{Ga}_{0.67}\text{As}$ spacer, a 38 nm n -doped $\text{Al}_{0.33}\text{Ga}_{0.67}\text{As}$ layer and a 17 nm undoped GaAs caplayer.

First a wet-etched Hall bar with 12 ohmic Ni/AuGe/Ni contacts is made on this material by optical lithography. In the middle of this mesa a grid-shaped gate is fabricated. The grid structure, consisting of 3×3 cells, is defined in a double layer of PMMA (with a total thickness of $\approx 200 \text{ nm}$) by nanolithography with a 100 keV electron beam, to reduce negative influences of the proximity effect between the closely packed wires of the grid. After development the gate metallization is evaporated as 10 nm titanium and 20 nm gold, at a background pressure below 10^{-7} mbar . The result after lift-off in boiling acetone is shown in figure 3.1. A small separation is intentionally left open between the grid and the upper gate to avoid (electrostatic) potential differences across

Figure 3.1: *Micrograph of the metallic gate structure (light area); the white line represents $1\ \mu\text{m}$*

the grid.

The grid is connected by wide leads to bond pads which are situated outside the etched mesa. By applying a negative voltage to the lead the electron density below the gate can be reduced in a controlled way.

The width of the lines d is less than 50 nm, the size of the unit cell a is 300 nm. In this way a is comparable to the Fermi wavelength of the electrons. Moreover it is anticipated that phase coherence will be preserved over several cells, since in general the inelastic mean free path l_{in} is much larger than $l_e \gg a$.

3.3 Results of the measurements

For the measurements the samples are mounted in a dilution refrigerator. The differential resistances are measured by means of a current-biased ac lock-in technique at a temperature of about 40 mK. Always the voltage on the upper gate in figure 3.1 (without grid) is chosen such that the small 2DEG channel between this gate and the grid is depleted totally ($V_{\text{gate}} = -560\ \text{mV}$). In this way it is ensured that the electron transport

takes place underneath the grid, i.e. through the 2D periodic potential [4].

Figure 3.2: The differential conductance as a function of the magnetic field for several voltages on the grid ($T = 0.05$ K, $V_{\text{gate}} = -560$ mV; from top to bottom $V_{\text{grid}} = -360$ mV, -365 mV, -370 mV, -375 mV, -380 mV, -385 mV, -390 mV, -395 mV); the letters A, B and C refer to the corresponding orbits in figure 3.4

In figure 3.2 the data of the conductance measurements in low magnetic fields are presented for several voltages V_{grid} on the grid. The applied excitation current is 0.1 nA. Especially for not too negative grid voltages minima in the conductance are observed at $B \approx 0.18$ T (A), $B \approx 0.48$ T (B) and $B \approx 0.80$ T (C). Superimposed on these minima small oscillations can be seen which are periodic in B . The result of increasing the temperature is shown in figure 3.3: while the oscillations are strongly suppressed at a temperature below 1 K, the conductance minima are only weakly affected up to a temperature of 2 K. This sug-

Figure 3.3: The differential conductance as a function of the magnetic field at several temperatures ($V_{\text{grid}} = -385$ mV, $V_{\text{gate}} = -560$ mV; from top to bottom $T = 0.05$ K, 0.1 K, 0.2 K, 0.5 K, 1.0 K, 2.0 K); the upper curves have been offset for clarity, the letters A, B and C again refer to the orbits in figure 3.4

gests that the minima are due to a classical phenomenon, in contrary to the oscillations which have a quantummechanical origin.

3.3.1 Localized orbits

The radii R_c of the cyclotron orbits [5] which belong to the magnetic fields where the minima occur are respectively 330 nm, 125 nm and 75 nm. In figure 3.4 these cyclotron orbits are drawn on scale in the grid. It is rather striking that the largest orbit A just encloses four line crossings, the middle orbit B encircles one crossing and that the smallest orbit C just fits in one unit cell. The cyclotron orbits cross the grid lines

Figure 3.4: *Schematic view of the grid with the localized electron orbits corresponding to the conductance minima A, B and C (on scale)*

at positions where the potential barrier is expected to be smallest. It is clear that at a certain V_{grid} the electron density under the crossings (where the total gate area in the vicinity is larger) will be lower than that under the lines connecting them [6]. Therefore at not too large values of V_{grid} a kind of antidot array is formed by the line crossings. Conductance minima due to commensurability of cyclotron orbits and an antidot array has earlier been reported [7, 8], which makes this explanation of the minima in our measurements very plausible. Normally the conduction through such an array is carried by electrons scattering between the antidots. However, at those magnetic field values where the cyclotron orbit “fits” in the array, a fraction of the conduction electrons dwells in these orbits, which will reduce the conductance through the array, yielding a conductance minimum. Clearly, this is only a largely simplified (but appealing) model. In reality the shape of the cyclotron orbits will deviate from a circle because of the varying two-dimensional potential, but this will not affect the validity of the basic idea of localized orbits. In fact this phenomenon is the two-dimensional variant of the Weiss oscillations [1] (first found in a one-dimensional periodic potential), which are theoretically described (semi-classically!) in [9].

3.3.2 Aharonov-Bohm oscillations

As is evident from the figures 3.2 and 3.3 clear oscillation patterns are found in the minima of the conductance. Nihey et al. [7, 10] also found oscillations in a single conductance minimum. The authors contribute these oscillations to the Aharonov-Bohm effect [11] in one cyclotron orbit. We find for the first time oscillations in three conductance minima. The period of the Aharonov-Bohm oscillations associated with the areas of the localized orbits are 12 mT, 86 mT and 230 mT. The measured periods are, however, ≈ 12 mT, ≈ 21 mT and ≈ 24 mT. So, while the period corresponding to the largest localized orbit A is in good agreement with the calculation, there is a discrepancy for the smaller orbits B and C. The difference between the calculated and the measured period is respectively about a factor 4 (B) and a factor 9 (C). This is exactly the number of localized orbits of kind B or C which is possible within the grid! The largest orbit A can be positioned only in one way in the grid, for the smaller ones there are however 4 and 9 possibilities. This leads us to the assumption that the localized orbits are phase-coherently coupled throughout the grid. This is reasonable, since in general the phase coherence length is larger than l_e and in the plain 2DEG the elastic mean free path $l_e \approx 6 \mu\text{m}$ is much larger than the total grid (which is 950 nm). The 4-fold, respectively 9-fold degeneracy of the localized states in case B and C is lifted by quantummechanical tunneling between these states. This leads to a splitting of the Aharonov-Bohm oscillations and thus to a reduction of the period by a factor 4, respectively 9. With this assumption taken into account the theoretically expected periods are 12 mT, 21 mT and 26 mT, which is in much better agreement with the experimental results.

3.4 Discussion of the measurements

From 4-terminal measurements of the resistance at very low magnetic fields (i.e. below 0.2 T) indications for weak localization effects have been found. Because of the 4-terminal configuration these effects cannot be due to the ohmic contacts (current and voltage probes are much further apart than the phase coherence length), so they must be attributed to the grid structure [12]. From a measurement for $V_{\text{grid}} = -330$ mV and $V_{\text{gate}} = -560$ mV, which shows a resistance minimum at $B = 0.12$ T, the phase coherence length is roughly estimated to be $\approx 1.6 \mu\text{m}$ (assuming

the pure metallic regime and the width of the relevant constriction being 250 nm) [13]. This confirms that phase coherence is preserved over the total size of the grid, even if localization due to the constriction in the grid is taken into account.

The oscillations can not be explained by the Shubnikov-De Haas effect, because they are periodic in B instead of in $1/B$. They are also fundamentally different from the magnetoconductance oscillations observed in single quantum dots in high magnetic fields [14], since our measurements are clearly not in the edge channel regime. In [8, 15, 16] a more rigorous theoretical model (quantization of chaotic electron motion) is presented, which takes into account the deviations from a hard wall potential. However, the much simpler “pinball” model with free cyclotron radii (described above) seems to be sufficient to give a good description of our measurements.

To make another consistency check we estimate the relevant energy scale of the oscillations. The total magnetic flux enclosed by the coupled orbits in minimum B at $B = 0.48$ T approximately equals 22 flux quanta ϕ_0 . With a Fermi energy $E_F = 4.7$ meV each flux quantum ϕ_0 corresponds to an energy of ≈ 0.21 meV. From this argument one predicts that the oscillations will be suppressed when $4k_B T$ is equal to 0.21 meV, i.e. at $T \approx 0.6$ K. Analogously we estimate that the oscillations in the minima A and C will be suppressed at $T \approx 0.9$ K and $T \approx 0.5$ K. This is in good agreement with figure 3.3. The results of the foregoing are identical to those following from the argument used by Nihey et al. to estimate the typical energy splitting [7].

The conductance minima at magnetic fields corresponding to localized orbits have also been observed in a similar grid with 9×15 unit cells (also with $a = 300$ nm and $d \approx 50$ nm). In this case no Aharonov-Bohm oscillations have been observed. This is clear, since now the total size of the grid ($4.55 \mu\text{m}$) is much larger and phase coherence all over the grid is not very likely.

If we look at the effect of variation of V_{grid} in figure 3.2, we see that the position of the conductance minima nearly stays unchanged, but that their amplitude decreases strongly at more negative V_{grid} . This indicates that the height of the potential barriers under the grid lines is increased rather than the width.

That the oscillations stay clearly visible shows again that this phenomenon has a different origin. Not a conductance minimum, but localization and phase coherence are required. When the potential land-

scape is firmly formed, so that the areas of low barrier height between the antidots are smallest, the cyclotron orbits are best localized, which is optimal for the Aharonov-Bohm oscillations. In fact this requirement is analogous to the condition of small channel width in the usual ring geometry for Aharonov-Bohm measurements.

We believe that Coulomb effects are not important for the description of our measurements. If we measure the magnetoconductance at different values of V_{gate} (varied from -480 mV to -580 mV), the oscillations in the minima are unchanged. Moreover, when V_{grid} is swept, while the magnetic field is fixed at the value of a conductance minimum, no oscillations are found in the conductance.

3.5 Summary

In summary, we have experimentally studied electron transport in a two-dimensional superlattice. Conductance minima due to localized (cyclotron) orbits have been measured. Moreover, for the first time Aharonov-Bohm oscillations superimposed on these minima have been observed that give evidence for phase-coherent, two-dimensional coupling of the localized orbits throughout the grid.

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Chapter 4

Indications for band-structure effects in two-dimensional lateral superlattices

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Abstract

Band structure effects have been investigated in two-dimensional lateral superlattices, consisting of 3×3 , 7×7 and 9×15 cells, which form tuneable arrays of quantum dots. In the quantum Hall regime oscillations in the longitudinal conductance have been found, which are periodic in the magnetic field. The oscillation period scales with the total number of

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dots in the array, in correspondence with the behaviour of the superlattice as an artificial two-dimensional crystal. The order of magnitude of the period is in very good agreement with this description. On the slope between the Hall plateaux a more irregular behaviour of the magnetoconductance has been measured, which can be attributed to the crossing of energy gaps by the Fermi level.

4.1 Introduction

Since the suggestion by Esaki and Tsu [1] that a superlattice will introduce minibands and minigaps, the interest in artificial band structures has been non-ceasing [2]. When a superlattice is made by epitaxial growth of alternating layers of different semiconductors, the potential modulation is one-dimensional. Lateral superlattices however provide the opportunity to make two-dimensionally periodically modulated potentials. Moreover, lateral superlattices formed by gate structures also give the possibility to tune the strength of the modulation potential.

Theoretically it has been predicted that the two-dimensional band structure becomes very exotic, when the magnetic length $l_m = \sqrt{\hbar/eB}$ is comparable to the lattice period a (“Hofstadter’s butterfly” [3]). In natural crystals magnetic fields of the order 10^4 T would be necessary to get into this regime. However, modern nanolithographic techniques recently made it possible to produce periodic semiconductor structures, which are suitable to study artificial band-structure effects at modest magnetic fields. In this paper we will describe magnetotransport measurements on several two-dimensional lateral superlattices in high magnetic fields where $l_m \lesssim a$ [4].

The samples are based on the two-dimensional electron gas (2DEG) in a GaAs/AlGaAs heterostructure, because of its high mobility, the large Fermi wavelength and its excellent gating properties. The used 2DEG has a mobility $\mu_e \gtrsim 100$ m²/Vs and an electron density $n_s \approx 1.3 \cdot 10^{15}$ m⁻², yielding a Fermi wavelength $\lambda_F \gtrsim 70$ nm and an elastic mean free path $l_e \gtrsim 6$ μ m.

On this material a grid-shaped gate is fabricated, consisting of 3×3 , 7×7 or 9×15 cells. (The fabrication process is described in more detail elsewhere [4].) The width of the lines d is approximately 50 nm, the size of the unit cell a is 300 nm (see figure 4.1). In this way a is comparable to the Fermi wavelength of the electrons. It is anticipated that phase coherence will be preserved over many cells, since in general the phase

Figure 4.1: *Micrograph of the metallic gate structure (light area) of the 7×7 grid; the lattice constant $a = 300 \text{ nm}$*

coherence length is much larger than the elastic mean free path and $l_e \gg a$.

By applying a negative voltage to the grid the electron density below the gate is reduced and a two-dimensional array of coupled quantum dots is formed. The strength of the modulation potential can be changed in a controlled way by varying the voltage on the gate. This gives an additional degree of freedom in the measurements, besides variation of the magnetic field.

For the measurements the samples are mounted in a dilution refrigerator. The differential resistances are measured by means of a current-biased ac lock-in technique at a temperature of about 40 mK. Always the voltage on the upper gate in figure 4.1 (without grid) is chosen such that the electron transport takes place underneath the grid, i.e. through the 2D periodic potential and not through the small 2DEG channel between this gate and the grid.

4.2 Results and discussion of the measurements

In figure 4.2 the magnetoconductance through the 7×7 -grid around the second Hall plateau at $G = 2e^2/h$ is shown for several voltages

Figure 4.2: *The differential conductance on the second Hall plateau as a function of the magnetic field for several voltages on the grid (from top down $V_{\text{grid}} = -172 \text{ mV}$, -178 mV , -184 mV , -190 mV ; $V_{\text{gate}} = -280 \text{ mV}$; the lower three curves have offsets in both directions); the inset shows the curves on a larger scale (without offsets)*

V_{grid} on the grid. The applied excitation current is 0.1 nA . Clear oscillations can be seen. To check the periodicity of these oscillations a Fast Fourier transformation has been performed after subtraction of the varying background. This background conductance has been determined by means of a Moving Average filter, which removed frequencies below 100 T^{-1} from the curves. The resulting Fourier spectra are presented in figure 4.3 which show clear peaks near 600 T^{-1} , corresponding to a period $\Delta B \approx 1.7 \text{ mT}$. In the following we will illustrate that these oscillations can be related directly to discrete states in the band structure of the 2D quantum dot array.

In the independent electron approximation the energy spectrum of a single quantum dot with radius R in a perpendicular magnetic field B

Figure 4.3: *Fourier spectra of the magnetoconductance shown in figure 4.2 (the upper curves have been offset for clarity); the B -interval over which the transformation has been done is shown in brackets*

can be described [5] by

$$E_{n,l} = \frac{1}{2}\hbar\omega_c l + (2n + |l| - 1)\sqrt{\frac{1}{4}\hbar^2\omega_c^2 + \hbar^2\omega_0^2} \quad (4.1)$$

(where the quantum number $n = 1, 2, 3, \dots$ and $l = 0, \pm 1, \pm 2, \dots$ and the cyclotron frequency $\omega_c = eB/m^*$), assuming a rotationally symmetric confining potential defined by

$$V(r) = \frac{1}{2}m^*\omega_0^2 r^2 \text{ with } \omega_0 = \sqrt{\frac{2E_F}{m^*R^2}}. \quad (4.2)$$

Using these equations we can estimate the typical splitting between the discrete energy states ΔE of a two-dimensional array consisting of q dots by calculating $(E_{n,l-1} - E_{n,l})/q$. For $R = 125$ nm ($\approx (a - d)/2$), $q = 49$ and $B = 1.75$ T we find $\Delta E = 4.3$ μ eV.

To check this experimentally the excitation current I has been varied. As can be seen in figure 4.4 the oscillations are suppressed if $I \gtrsim 0.33$ nA.

For a conductance of $2e^2/h$ this corresponds to an energy splitting of

Figure 4.4: *Influence of the excitation current on the magnetoconductance oscillations (the lower three curves have been offset vertically)*

4.2 μeV , which is in remarkable agreement with the theoretical prediction.

It is also possible (and in some aspects simpler) to describe the electron transport through the array by means of edge channels. This is allowed, since the classical cyclotron radius $R_c \lesssim 35$ nm is much smaller than the lattice period a . It is known that at high magnetic fields a new kind of the Aharonov-Bohm effect can be observed in a quantum dot, due to resonant tunneling and reflection of edge channels [2, 6]. The period in the magnetic field is given by $\Delta B \simeq h/eS$, with S the area enclosed by the circulating edge state (in fact a “zero-dimensional state”). In one quantum dot of our superlattice this would be ≈ 84 mT. However, if the 49 dots of the array are coupled, we have to divide this value by 49, which results in $\Delta B = 1.7$ mT. This is in remarkably good agreement with the measured period.

The Fourier spectra in figure 4.3 show a weak dependence of the oscillation period on V_{grid} : at more negative V_{grid} the period is slightly smaller. It could be that the effective area of the dots is larger, yielding a smaller period, as follows. If V_{grid} becomes more negative, the Hall plateaux with oscillations are observed at smaller magnetic fields and at

smaller magnetic fields the edge channel is situated closer to the edge, thus encircling a larger area. If this is true, it means that V_{grid} barely affects the dimensions of the dots, but only the coupling between them.

As long as the edge channel belonging to the second (spin-split) Landau level is well transmitted the conductance is $\approx 2e^2/h$ and oscillations due to the discrete states in the two-dimensional band structure can be observed. When the transmission of the channel however decreases (with increasing magnetic field), it is expected that at a certain point the energy gap, which opens for transmission probabilities smaller than 1, will be encountered [8,9]. As has been calculated by Kirczenow, the spectrum of energy levels is very complicated near this gap. Indeed B -intervals where the magnetoconductance behaves more or less irregularly have been measured on the step between the Hall plateaux (see inset of figure 4.2). We believe that in these regions the Fermi level crosses the gap region in the band structure. This agrees with the description of the quantum dot array as an artificial crystal.

Note that this behaviour is different from that of a single quantum dot: in that case the Aharonov-Bohm oscillations are found between two Hall plateaux, while the effect is suppressed on the plateaux themselves [6]. In our superlattices the periodic oscillations are found on the plateaux and not in between.

When the conductance has arrived at the first Hall plateau ($G \approx e^2/h$), the transmission is again (nearly) equal to 1, but now for the edge channel with the lowest Landau level index. So it is expected that again magnetoconductance oscillations will occur. This is indeed the case, as can be seen in figure 4.5. The period ΔB is the same as that on the second Hall plateau. At first sight this is surprising, since the magnetic field is now much larger. However, the value of the period is defined by the dimensions of the grid, which are fixed [7]. The invariability of the oscillations shows that they are due to the artificial crystal structure and not to simple Landau quantization or commensurability effects depending on magnetic field (like Shubnikov-de Haas or Weiss oscillations).

To make sure that this phenomenon is a two-dimensional effect, we also measured the magnetoconductance of similar grids with a different number of cells, respectively 3×3 and 9×15 . Again regions of complex behaviour are observed between the Hall plateaux, attributed to opening of an energy gap. In the 3×3 -grid regular oscillations on the first Hall plateau are found, as can be seen in figure 4.6. The oscillation period

Figure 4.5: *The differential conductance on the first Hall plateau as a function of the magnetic field (from top down $V_{\text{grid}} = -162$ mV, -174 mV, -186 mV, -198 mV; $V_{\text{gate}} = -280$ mV; the lower three curves have been offset in both directions)*

is determined by Fourier transformation to be $\Delta B \approx 8.7$ mT. Most important is the observation that this value is approximately $49/9 \times$ the period measured in the 7×7 -grid: this proves that the oscillations have a two-dimensional origin, because ΔB is inversely proportional to the total number of cells in the grid! Evidently, the value is also in agreement with the calculation based on the Aharonov-Bohm effect of edge channels, which predicts $\Delta B = 9$ mT.

In the 9×15 -grid noisy structures are measured on the second Hall plateau; the Fourier spectra do not show clear peaks. This is understandable, because the expected energy splitting is $1.5 \mu\text{eV}$ and the thermal energy $k_B T$ is already $3.4 \mu\text{eV}$ at $T = 40$ mK. Nevertheless the irregular behaviour which is observed between the Hall plateaux may be an indication for an effect caused by the two-dimensional potential.

Moreover, in the 3×3 -grid the conductance as a function of V_{grid} (at constant B) shows a structure that is comparable to the structure in the magnetoconductance, i. e. oscillations on the first Hall plateau and complex aperiodic behaviour between the first and the second plateau.

Figure 4.6: *The differential conductance of the 3×3 -grid as a function of the magnetic field (from top down $V_{\text{grid}} = -260 \text{ mV}$, -290 mV , -320 mV , -350 mV ; $V_{\text{gate}} = -600 \text{ mV}$)*

Careful measurements of the magnetoconductance at slightly different V_{grid} indicate that the magnetoconductance oscillations are shifted over one period, when V_{grid} changes 6 mV. This is also the typical period of the oscillations as a function of V_{grid} , which indicates that both oscillations have the same origin.

In measurements of oscillations like these one should be careful not to disregard Coulomb effects [10]. Nevertheless we think that Coulomb effects are not important for the interpretation of our results. In the first place we have observed oscillations at conductances above the critical value of $2e^2/h$ (see, for example, figure 4.2). Besides that, since it is known from other measurements that the spin degeneracy is already lifted at magnetic fields larger than $\approx 0.5 \text{ T}$, simple Coulomb effects are only expected to occur at conductances smaller than e^2/h . Also it is not clear why the period of oscillations would be inversely proportional to the total number of quantum dots in the superlattice, if they were due to Coulomb effects. The combination of parallel and in series coupling of dot capacitances suggests a different dependence.

4.3 Summary

In summary, for the first time (to our knowledge) experimental indications have been found for band-structure effects in two-dimensional superlattices. Oscillations in the magnetoconductance have been measured with a period that scales inversely with the total number of cells in the grid. The magnitude of the period is in very good agreement with theoretical calculations. In the 3×3 -grid the oscillations due to the band structure also have been observed as a function of the voltage on the grid-shaped gate. Moreover, indications for the existence of energy gaps have been found. We can conclude that the superlattice can be described as an artificial two-dimensional crystal.

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Chapter 5

Theory of interfaces between superconductors and normal materials

5.1 Introduction

The combination of superconductors and semiconductor nanostructures forms a challenge for the modern solid state physicists and the rapid developments in submicron technology makes it experimentally feasible. For a fruitful theoretical description of these systems also the accepted concepts of both research fields have to be combined. This integration process is evolving, but not yet finished.

In superconductivity BCS-like models [1] play a central role, while in mesoscopic semiconductor systems the Landauer-Büttiker formalism [2, 3] has proven its value. While in superconductors it is not “done” to ask which exact pathes a specific, individual particle can follow, mesoscopic semiconductor structures are described by transmission and reflection probabilities of single particles. This often allows to divide complex mesoscopic systems in subsystems which can be modeled by means of a relatively simple scattering matrix.

Since many papers have been and are written on the theory of junctions of superconductors (S) and normal materials (N), this chapter will be limited to the introduction of some basic concepts that are presently used in literature, in order to provide some background for the following chapters of this thesis. First the more traditional way of thinking about NS interfaces will be treated: in terms of the proximity effect. After that a special reflection process at NS-interfaces, Andreev reflection, is de-

scribed. This phenomenon recently got renewed attention, since it may provide a very elegant way to incorporate superconducting boundaries in the description of mesoscopic structures. Some theoretical models based on Andreev reflection will be discussed.

5.2 The proximity effect

In experiments it has been observed that a thin superconducting film deposited on a normal metal has a transition temperature which may be significantly lower than in bulk superconducting material; this is not the case for superconducting samples on nonmetallic substrates. Other experiments have shown that a thin ($\sim 10^2$ nm) normal slab sandwiched by two superconductors S and S' is able to carry a finite supercurrent from S to S'. De Gennes [4,5] first described a model, which is now known as the proximity effect, to explain these experimental observations.

He proposes that Cooper pairs can diffuse into a normal material, if there is a good electrical contact between the N and S regions. Algebraically this means that the pair potential $\Delta(r)$ extends up to a rather long distance, with characteristic length ξ_N , into the N region.

The pair potential is defined by the wellknown BCS self-consistency equation

$$\Delta(r) \equiv V(r) \langle \psi_{\uparrow}(r) \psi_{\downarrow}(r) \rangle, \quad (5.1)$$

where $\psi_{\uparrow}(r)$ and $\psi_{\downarrow}(r)$ are one-electron operators, $V(r)$ is the strength of the electron-electron interaction and the brackets denote a thermal average. The quantity $V(r)$ will not be the same in the N and S regions. In the S regions $V = V_S$ is positive (attractive interaction), while in the N regions $V = V_N$ is usually negative for semiconductors (repulsive interaction)¹. Following De Gennes [4], the pair potential $\Delta(r)$ is a natural “order parameter” for inhomogeneous systems. This is not the only possible choice; an other candidate is the “condensation amplitude” $F(r)$ defined by

$$F(r) = \langle \psi_{\uparrow}(r) \psi_{\downarrow}(r) \rangle. \quad (5.2)$$

F is the probability amplitude of finding two electrons in the condensed state at point r . Although both $\Delta(r)$ and $F(r)$ are continuous functions

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¹It is also possible that $V_N > 0$, if the N material is a superconductor at low enough temperatures. In general, the sign of V_N depends on a delicate balance between Coulomb repulsion and phonon-induced attraction.

of r on an atomic scale, on a larger scale neither of them is continuous at the interface between the two metals, if this interface is modeled as a sharp boundary. For dirty systems, where the transport is diffusive, the quantity $\frac{F(r)}{N(r)} = \frac{\Delta(r)}{N(r)V(r)}$ is continuous, where $N(r)$ is the local density of states (per energy unit and per volume unit) at the Fermi level. The results of calculations of an NS boundary can qualitatively be represented in the N region ($x > 0$) in the form

$$\Delta(x) \sim \Delta_0 \frac{N_N}{N} T_j e^{-\frac{x}{\xi_N}}, \quad (5.3)$$

where $\frac{N_N}{N}$ is the ratio of the local densities of states at the Fermi surface in the N respectively the S region, and T_j is the transmission coefficient at the boundary for an electron at the Fermi surface. This dependence of the order parameter is shown in figure 5.1. The characteristic diffusion

Figure 5.1: Variation of the order parameter at an NS interface

length is for dirty metals given by

$$\xi_{N,S} = \sqrt{\frac{\hbar D_{N,S}}{2\pi k_B T}}, \quad (5.4)$$

with $D_{N,S}$ the diffusion constant. The distance over which the order parameter in the S region is suppressed (characterized by ξ_S) is, in general, very small compared to ξ_N ; therefore this reduction is often ignored (rigid boundary conditions).

Figure 5.2: Variation of the order parameter in an SNS junction

In an SNS junction finite order parameters are induced in the N region from both superconductors. This is schematically shown in figure 5.2. If the order parameter is non-zero throughout the normal region, a supercurrent can flow from one superconductor to the other.

Often $\Delta(r)$ is confused with the energy gap Δ_0 ; the latter is the lowest excitation energy of the fermions in the self-consistent field Δ and is independent of r .

If semiconductors are used for the N region, a jump in the order parameter will occur at an NS boundary because of the Schottky barrier. Therefore lower induced order parameters are expected in semiconductors, compared to metals. For semiconductors² it is convenient to express formula (5.4) as

$$\xi_N = \sqrt{\frac{\hbar^3 \mu}{6\pi e m^* k_B T}} \sqrt[3]{3\pi^2 n} \quad (5.5)$$

for the three-dimensional case, and

$$\xi_N = \sqrt{\frac{\hbar^3 \mu n}{2e m^* k_B T}} \quad (5.6)$$

for the two-dimensional case [7]. In these formulas μ is the electron mobility, n the electron density and m^* the effective mass of an electron.

²These expressions are correct for n -type III-V compounds; a more general formula has been derived in Ref. [6].

For a two-dimensional electron gas in GaAs/AlGaAs heterostructures (with typically $\mu = 10^2 \text{ m}^2/\text{Vs}$ and $n = 2 \times 10^{15} \text{ m}^{-2}$) a value for ξ_N of circa $4 \text{ }\mu\text{m}$ is expected at $T = 50 \text{ mK}$.

5.3 Andreev reflection

Besides the rather phenomenological proximity-effect model, microscopic theories based on scattering processes have arisen. In 1964 Andreev [8] described a new kind of reflection process which can occur at an NS interface. This phenomenon is nowadays known as Andreev reflection.

Consider an electron in the N region of a NS sandwich, with an excitation energy E below the (bulk) energy gap of the superconductor Δ_0 . This electron cannot penetrate in the superconducting region, so it has to undergo reflection. As described in Ref. [9] the momentum p of the electron can have a value close to the Fermi momentum p_F , while the distance over which the energy gap changes from zero in the N region to the bulk value in the S region equals the BCS coherence length $\xi_0 \equiv \frac{\hbar v_F}{\pi \Delta_0}$. The corresponding uncertainty in the momentum

$$\delta p \sim \frac{\hbar}{\xi_0} = \frac{\pi \Delta_0}{2E_F} \cdot p_F, \quad (5.7)$$

which is much smaller than p_F . Thus the inhomogeneity is too weak for ordinary reflection processes that change p by a quantity of the order of p_F . It is, however, possible that the electron is reflected as a hole (with excitation energy $-E$), which is called Andreev reflection. While normal reflection only changes the sign of the velocity component normal to the boundary, in Andreev reflection all velocity components change sign (figure 5.3). In this case the excitation momentum is practically unchanged. The laws of charge, energy and particle conservation are satisfied, since a Cooper pair is added to the superconductor.

Although it is often convenient to describe Andreev reflection as a momentary process taking place at a sharp NS boundary, it can also be thought to happen more gradually, as shown in figure 5.4 [10]. The local excitation spectrum changes with position, because the energy gap gradually increases. While the momentum of the initial electron-like excitation A changes, its energy stays constant. At the turning point, where this energy is equal to the gap, A'' coincides with B''. This state emerges back in the normal region to a hole-like excitation B. In this

Figure 5.3: *Normal and Andreev reflection in real space*

Figure 5.4: *The excitation spectra at several positions on a gradual NS boundary*

way Andreev reflection can be described as an adiabatic transition from an electron-like to a hole-like state.

In many calculations dealing with Andreev reflection the (time-independent) Bogolyubov-de Gennes equations [5, 11] take up a central place:

$$\mathcal{H}\Psi = E\Psi. \quad (5.8)$$

This formula has the same form as the Schrödinger equation; however,

Ψ is a two-component wave function

$$\Psi = \begin{pmatrix} u(x) \\ v(x) \end{pmatrix} \quad (5.9)$$

and the Hamiltonian $\mathcal{H}$ is a matrix

$$\mathcal{H} = \begin{pmatrix} H - E_F & \Delta \\ \Delta^* & -(H - E_F)^* \end{pmatrix}, \quad (5.10)$$

where $H = -\frac{\hbar^2}{2m}\nabla^2$ is the kinetic energy operator. In these equations $u(x)$ describes the electron-like wave function and $v(x)$ the hole-like wave function with an excitation energy E relative to the Fermi energy E_F . If $\Delta \equiv 0$ the equations for both wave functions are decoupled and become normal Schrödinger equations. The strength of the Bogolyubov-de Gennes equations is the possibility to describe spatially inhomogeneous systems (such as SNS systems) in a insightful way, since it is often convenient to think in terms of electrons and holes. For example, they are very suitable for the calculation of transmission and reflection probabilities of excitations in SN(S) systems. Hereby a good starting-point is provided for the description of mesoscopic superconductor/semiconductor systems. This is illustrated in the next chapter, in which a transmission formalism is used to calculate supercurrent flow in such a system.

An important application of the Bogolyubov-de Gennes equations is the theory for the $I - V$ -curves of NS microconstriction contacts, developed by Blonder, Tinkham and Klapwijk [12]. Using a generalized semiconductor model, probabilities for normal and Andreev reflection have been calculated for an NS interface. At the interface a delta-function shaped potential is located acting as a barrier. The strength of this barrier is expressed in the dimensionless parameter Z , which is related to the transmission coefficient T of the barrier by

$$T = \frac{1}{1 + Z^2}. \quad (5.11)$$

By varying this parameter the range from a clean metallic contact ($Z = 0$) to a tunnel junction ($Z \gtrsim 10$) can be covered. Much valuable information about the $I - V$ -curves of an NS boundary is obtained from the calculated transmission and reflection probabilities. For example, it is found that the subgap resistance is smaller than the normal-state resistance for $Z \lesssim 0.5$, while it grows rapidly with increasing Z . This method has also been applied to an SNS system, which yielded an explanation

for the observed subharmonic gap structure in terms of multiple Andreev reflection [13]. However, because all phase information is ignored in the calculation, no interference effects are included.

Andreev [9] already realized that the periodic motion of particles caused by the reflections in a phase-coherent SNS system can lead to a quantization of energy levels below the superconducting gap. This has been worked out by Kulik [14]. He makes the intriguing statement: "... this quantization implies the presence of a coherent connection between the phases of the ordering parameters of the superconductors adjacent to the normal material, i. e. the existence of superconducting ordering, for the entire SNS system as a whole. (...) On the other hand, phase coherence denotes the possible flow of superconducting currents (the Josephson effect) through the junction." This is probably the first instance where Andreev reflection is associated with supercurrent flow through a normal material.

By applying the Bogolyubov-de Gennes equations to a (quasi one-dimensional) SNS system as shown in figure 5.5 ($\Delta(x) = 0$ in the N

Figure 5.5: Quantization of energy levels of excitations in a SNS system; the dashed lines represent $\phi_1 = \phi_2$, the thin solid lines $\phi \neq \phi_2$

region, $\Delta(x) = \Delta_0 e^{i\phi_{1,2}}$ in superconductor 1,2) the following expression is obtained for the quantized energy levels [14–16]

$$E_n^\pm = \frac{\hbar v_F}{2L^*} \left[2\pi n + 2 \arccos \left(\frac{E}{\Delta_0} \right) \mp (\phi_1 - \phi_2) \right], \quad (5.12)$$

with $n = 0, 1, 2, \dots$ and L^* the effective length between the superconductors [15]. E_n^+ are energy states of electrons traveling in positive x direction and holes in negative x direction; E_n^- denote energy states

with the directions reversed. When $\phi_1 = \phi_2$ the states are degenerate ($E_n^+ = E_n^-$). From this energy spectrum the supercurrent can be calculated for a given phase difference as described in the next chapter. This means that phase-coherent Andreev reflection and supercurrent flow are, at least, very closely related.

Recently, new interference effects caused by Andreev reflection have been theoretically studied [17–19]. Some of them will be discussed together with our experiments in chapter 9 and 10. While all foregoing theories studied in fact ballistic systems, now also elastic scattering is considered, which of course doesn't destroy phase coherence. An important result is the general expression for the conductance of an NS junction in terms of transmission coefficients [18]:

$$G_{\text{NS}} = \frac{4e^2}{h} \sum_{n=1}^N \frac{T_n^2}{(2 - T_n)^2}, \quad (5.13)$$

where T_n ($n = 1, 2, \dots, N$) are the eigenvalues of the Hermitian transmission matrix. The power of this formula, which is the multichannel generalization of a formula for the single channel case in Ref. [12], is comparable to that of the Landauer formula for the normal-metal conductance:

$$G_{\text{N}} = \frac{2e^2}{h} \sum_{n=1}^N T_n. \quad (5.14)$$

5.4 Concluding remarks

It is clear that both formalisms have yielded valuable insights and results. While originally the proximity effect was intended for diffusive systems and the description in terms of Andreev reflection for ballistic systems, both theories are now growing towards each other. Evidently, Andreev reflection can have a large influence on interference effects like weak localization, as will be shown in detail in chapters 9 and 10.

Especially now phase-coherent Andreev reflection has been associated theoretically with supercurrent flow, the question arises whether there exists any principal difference between a description in terms of electrons and holes, and the proximity-effect model. For example, the expression for the diffusion length of Cooper pairs, equation (5.4), is similar to the definition of the thermal length $l_{\text{T}} \equiv \sqrt{\frac{\hbar D}{k_{\text{B}} T}}$, used for electrons in normal diffusive systems.

If we naively try to make an estimation for the maximum $I_c R_N$ product of the system in figure 5.5 by substituting $R_N = h/2e^2$ (for a one-dimensional channel) and $I_c = \frac{ev_F}{L^*}$ (see next chapter) in the well-known relation

$$I_c R_N \sim \frac{2\Delta}{e}, \quad (5.15)$$

we find that the relevant superconducting gap has to be $\Delta = \frac{\hbar v_F}{4L^*}$. It is difficult to believe that this is accidentally equal to the value of $E_0(\phi_1 = \phi_2)$ in formula (5.12).

More theoretical and certainly also experimental work (with semiconductors as well as with normal metals) is desirable to clarify the correspondences between the proximity-effect model and models based on phase-coherent Andreev reflection. In our opinion this is important, since a complete theory developed around Andreev reflection seems to provide a powerful and elegant formalism, that is very suitable for the study of mesoscopic superconductor/semiconductor systems. Moreover, the last criterion in science is an aesthetic criterion.

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Chapter 6

Transmission formalism for supercurrent flow in multiprobe superconductor-semiconductor-superconductor devices

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Abstract

A theoretical study is given of supercurrent flow in a one-dimensional semiconductor channel coupled to superconductors at both ends. In addition, the channel is coupled to a semiconductor reservoir by means of a junction with variable coupling strength ϵ . The supercurrent $I(\phi)$ is calculated from the phase-coherent propagation of electronlike and holelike excitations emitted by the superconductor reservoirs, together with electron and hole excitations from the semiconductor reservoir. The effect of temperature and ϵ on $I(\phi)$ is studied. It is shown that a voltage applied between the semiconductor reservoir and the superconductors modifies the $I(\phi)$ relation, even in the limit $\epsilon \rightarrow 0$.

6.1 Introduction

In recent years our understanding of electron transport in mesoscopic conductors has greatly improved. It has become clear that at low temperatures electrons can maintain their phase coherence over considerable distances. Recently it has also become possible to study devices through which the electron can travel ballistically, without being scattered by impurities.

A challenge is now presented by the possibility to combine this ballistic, phase-coherent transport with superconductivity. A promising system is a high-mobility two-dimensional electron gas (2DEG), in which a narrow channel is defined (e.g., by means of a quantum point contact [1]), and connected to one or more superconducting electrodes. Very recently it was predicted [2] that the supercurrent in such a device, measured as a function of the width of the channel, should exhibit steps each time an additional one-dimensional channel is opened.

A large volume of literature exists on the theory of supercurrent flow through superconductor-normal-metal-superconductor (S - N - S) or superconductor-semiconductor-superconductor (S - SM - S) systems. A selection is given in Refs. [2–8]. In most of these papers the supercurrent is calculated using Green's-function methods [9]. In this paper we will describe the supercurrent flow through a S - SM - S device with the use of a transmission formalism. This formalism has already been applied for the description of electron transport in normal metals and semiconductors [10, 11]. It has also been used for the description of transport through superconductor-normal-metal interfaces [12–14].

6.2 Description of the formalism

The device geometry is illustrated in Fig. 6.1. Leads 1 and 2 (with lengths L_1 and L_2) form a one-dimensional (1D) channel [15], which is connected at both ends to superconductors 1 and 2 (sc1 and sc2), with electrochemical potentials μ_1 and μ_2 . By means of lead 3, the 1D channel can also be coupled to a wide 2D region, which can be described as a semiconductor reservoir with electrochemical potential μ_3 . The transmission and reflection of the electron waves at the junction formed by leads 1, 2 and 3, can be described by an $\underline{\mathbf{S}}$ matrix [16, 17], which relates the amplitudes of the outgoing waves to the amplitudes of the

Figure 6.1: *Layout of the system, which consists of a 1D channel formed by leads 1 and 2, coupled to superconductors 1 and 2. Lead 3 couples the channel to a semiconductor reservoir 3.*

incoming waves at the junction

$$\underline{\mathbf{S}} = \begin{bmatrix} -(a+b) & \epsilon^{1/2} & \epsilon^{1/2} \\ \epsilon^{1/2} & a & b \\ \epsilon^{1/2} & b & a \end{bmatrix} \quad (6.1)$$

with $a = 1/2[\sqrt{(1-2\epsilon)}-1]$, and $b = 1/2[\sqrt{(1-2\epsilon)}+1]$. The parameter ϵ describes the strength of the coupling between the 1D channel and the semiconductor reservoir. Maximum coupling is achieved for $\epsilon = \frac{1}{2}$, the channel is completely decoupled when $\epsilon = 0$.

We use a well-known model [2–6] where the pair potential $\Delta(x) = 0$ in the semiconductor channel, and $\Delta(x) = \Delta_0 \times \exp(i\phi_{1,2})$ in sc1 and sc2, with $\phi_{1,2}$ the superconductor phases. The wave functions can be found from the time-independent Bogolyubov-de Gennes equation [18]

$$E \begin{bmatrix} u(x) \\ v(x) \end{bmatrix} = \begin{bmatrix} H - E_F & \Delta \\ \Delta^* & -(H^* - E_F) \end{bmatrix} \begin{bmatrix} u(x) \\ v(x) \end{bmatrix} \quad (6.2)$$

In Eq. (6.2) $u(x)$ describes the electron wave function, and $v(x)$ describes the hole wave function with an excitation energy E relative to the Fermi energy E_F . There are four types of solutions. In the leads 1, 2, and 3,

$$\begin{bmatrix} u(x) \\ v(x) \end{bmatrix} = \begin{bmatrix} 1 \\ 0 \end{bmatrix} \exp(\pm iq^+ x), \text{ with } \hbar q^+ = \sqrt{2m}\sqrt{E_F + E} \quad (6.3a)$$

with $+$ ($-$) corresponding to electron excitations which move in the pos-

itive(negative) x direction [19], and

$$\begin{bmatrix} u(x) \\ v(x) \end{bmatrix} = \begin{bmatrix} 0 \\ 1 \end{bmatrix} \exp(\pm i q^- x), \text{ with } \hbar q^- = \sqrt{2m} \sqrt{E_F - E} \quad (6.3b)$$

with $+$ ($-$) corresponding to hole excitations which move in negative (positive) x direction. In the superconductors the four solutions are [20]

$$\begin{bmatrix} u(x) \\ v(x) \end{bmatrix} = \begin{bmatrix} u_0 \exp(i\phi) \\ v_0 \end{bmatrix} \exp(\pm i k^+ x),$$

with $\hbar k^+ = \sqrt{2m} \sqrt{E_F + (E^2 - \Delta_0^2)^{\frac{1}{2}}}$ (6.4a)

with $+$ ($-$) corresponding to electronlike excitations which travel in positive(negative) x direction, and

$$\begin{bmatrix} u(x) \\ v(x) \end{bmatrix} = \begin{bmatrix} v_0 \exp(i\phi) \\ u_0 \end{bmatrix} \exp(\pm i k^- x),$$

with $\hbar k^- = \sqrt{2m} \sqrt{E_F - (E^2 - \Delta_0^2)^{\frac{1}{2}}}$ (6.4b)

with $+$ ($-$) corresponding to holelike excitations which travel in negative(positive) x direction. The coherence factors are given by

$$u_0 = 1/\sqrt{2}\{1 + [1 - (\Delta_0/E)^2]^{\frac{1}{2}}\}^{\frac{1}{2}}, \quad (6.5a)$$

$$v_0 = 1/\sqrt{2}\{1 - [1 - (\Delta_0/E)^2]^{\frac{1}{2}}\}^{\frac{1}{2}}. \quad (6.5b)$$

The formalism we use is based on Büttiker's description of phase-coherent electron transport in multiprobe normal conductors [11]. The physical basis of this description is that at sufficiently low temperatures inelastic processes are negligible in leads 1, 2, and 3, and occur exclusively in the reservoirs. In normal conductors these reservoirs can be thought of as emitters (and absorbers) of electron waves. We extend this description by treating the superconductors as reservoirs also. A semiconductor reservoir can emit (and absorb) electron and hole excitations with energies $E > 0$. However, because no propagating solutions of Eq. (6.2) exist in the superconductor for $E < \Delta_0$, the superconductors can emit (and absorb) electronlike and holelike excitations for $E > \Delta_0$ only.

An important feature of 1D transport is that the product of one-dimensional density of states $N(E) = 2/\pi(dE/dk)^{-1}$ and the group velocity $v_g(E) = 1/\hbar(dE/dk)$ is independent of E and equal to $4/\hbar$ (including both spin directions). This result holds for electron and hole

states in the semiconductor, but also for the electronlike and holelike states in the superconductors [13]. It follows that the *particle* current [21] which is emitted by a reservoir per unit energy is equal to $2/h$.

Our description can also be used to describe normal current. Normal currents can be defined as currents which flow as a result of a difference between μ_1 , μ_2 , and μ_3 . In this paper we focus on the supercurrent which flows between sc1 and sc2, and its dependence on the phase difference $\phi = \phi_1 - \phi_2$. We define the reflection probabilities $R_{nn}(\phi, E)$ as the *charge* currents (in units of e) which flow back into the n th reservoir as a result of the *particle* current which is emitted by the n th reservoir at energy E .

Similarly the transmission probabilities $T_{nn}(\phi, E)$ are defined as the ratio of the charge currents which flow into the n th reservoir as a result of the particle currents emitted by the m th reservoir at energy E . Because the reservoirs can emit two types of particles, we add a superscript p , with $p = 1$ for electron (or electronlike) excitations, and $p = 2$ for hole (or holelike) excitations. We now define a charge current density (current per unit energy) $J_n^{ex}(\phi, E)$ in lead n by

$$J_n^{ex}(\phi, E) = \frac{2e}{h} \sum_{p=1,2} C_n^p(E) - R_{nn}^p(\phi, E) - \sum_{m \neq n} T_{nm}^p(\phi, E). \quad (6.6)$$

In Eq. (6.6) $C_n^p(E)$ indicates the ratio between the charge current (in units of e) and the particle current carried by the excitations which are emitted by the n th reservoir. $C_n^p(E) = -1$ for electrons, 1 for holes, $(v_0^2 - u_0^2)^{-1}$ for electronlike, and $(u_0^2 - v_0^2)^{-1}$ for holelike excitations. $J_n^{ex}(\phi, E)$ can be considered as the contribution to the supercurrent in lead n from excitations with energy E , under the condition that the occupation probability of the excited states in the reservoirs is *unity*.

We calculated the supercurrent $I_n(\phi, T)$ in lead n ($n=1,2$), with $\mu_1 = \mu_2 = \mu_3$. At $T = 0$ no excitations are present, and this seems to prevent the calculation of $I_n(\phi, T = 0)$. We can now perform a simple manipulation [22], and write the supercurrent at a finite T as

$$I_n(\phi, T) = I_n(\phi, T = 0) + I_n^{ex}(\phi, T), \quad (6.7)$$

with $I_n^{ex}(\phi, T)$ the supercurrent generated by the excitations at temperature T . Since the occupation of the excited states is given by $f(E, T) = 1/[1 + \exp(E/kT)]$, and $I_n(\phi, T)$ should vanish for $T \rightarrow \infty$, we can write [23]:

$$I_n(\phi, T) = \int_0^\infty [f(E, T) - \frac{1}{2}] J_n^{ex}(\phi, E) dE. \quad (6.8)$$

Equations (6.6) and (6.8) show that the supercurrent is expressed as an incoherent sum of the contributions of excitations at different energies, and from different reservoirs.

We have calculated the energy and phase dependence of $R_{nm}^p(\phi, E)$ and $T_{nm}^p(\phi, E)$ with an algorithm which matches the wave functions (6.3) and (6.4) at the superconductor-semiconductor interfaces, taking into account the phase shifts in leads 1 and 2, and matches the wave functions at the junction according to the $\underline{\mathbf{S}}$ matrix given in (6.1).

The fundamental mechanism which makes supercurrent flow possible is that of *coherent Andreev reflection* [24]. An electron in the 1D channel with $E < \Delta_0$ is reflected at the interface with sc2 as a hole, and its wave function can be shown to acquire a phase factor $\exp[i(-\phi_2 - \pi/2)]$. The reflected hole is converted back to an electron at the interface with sc1, and acquires an additional phase factor $\exp[i(\phi_1 - \pi/2)]$. This implies that when $\epsilon = 0$, bound states are formed in the semiconductor channel [3–6]. The energies of these states (positive states) depend on $(\phi = \phi_1 - \phi_2)$, and can be written as $E_n^+ = \hbar v_F / 2L^* [2\pi(n + \frac{1}{2}) + \phi]$, with $L^* = L_1 + L_2 + \hbar v_F / \Delta_0$, the effective length between the superconductors [25]. Similarly, the energy of the negative states, which correspond to electrons traveling in negative x direction (and holes traveling in positive x direction), can be written as

$$E_n^- = \hbar v_F / 2L^* [2\pi(n + \frac{1}{2}) - \phi]. \quad (6.9)$$

6.3 Results of the calculations

For the calculation we have taken [26] $E_F = 10\Delta$, $k_F L_1 = 77.9$, and $K_F L_2 = 94.9$. For a typical 2DEG in a GaAs/Al_xGa_{1-x}As heterostructure with effective mass $m^* = 0.067m_e$, and $E_F = 10$ meV, this corresponds to $\Delta = 1$ meV, and $L_1 + L_2 \approx 1.5$ μm . Figure 2 shows the calculated $J_1^{ex}(\phi, E)$ in lead 1 for $\epsilon = 0.2$. For $\phi \neq 0$ and $\phi \neq \pi$, $J_1^{ex}(\phi, E)$ for $E < \Delta_0$ consists of a series of positive and negative peaks, which correspond to the positive and negative states discussed above. For $\phi = 0$ the positive and negative states are all at the same energy, and $J_1^{ex}(\phi, E) = 0$. When ϕ is increased, the positive and negative states shift away from each other, until at $\phi = \pi$, the states again coincide, resulting in $J_1^{ex}(\phi, E) = 0$.

For $E < \Delta_0$, only excitations from the semiconductor reservoir contribute to $J_1^{ex}(\phi, E)$. Due to the coupling to this reservoir, the positive

and negative states are broadened in energy. When the coupling strength ϵ is reduced, the peaks develop into δ functions. It is interesting to note that when $\epsilon \rightarrow 0$, a finite supercurrent can be carried in an infinitesimally narrow energy interval. This should be contrasted with normal current flow between reservoirs, where the current per unit energy is limited to $2e/h$. The major contribution in the regime $E > \Delta_0$ is from the excitations from the superconductors [27]. For $E > \Delta_0$ there is only partial (Andreev) reflection at the S - SM interfaces. Consequently the positive and negative states are broadened substantially for $E > \Delta_0$ (even when $\epsilon = 0$). However, as can be seen in Fig. 6.2 these states can still give a (positive or negative) contribution to $J_1^{ex}(\phi, E)$.

Figure 6.2: Supercurrent density $J_1^{ex}(\phi, E)$ in lead 1 due to excitations with E , for a set of phase differences ϕ .

The supercurrent-phase relation $I_1(\phi, T)$, calculated from Eq. (6.8) is shown in Fig. 6.3, for several values of ϵ and T . For small values of ϵ , $I_1(\phi, T = 0)$ approaches a sawtooth shape (which was also obtained

Figure 6.3: Supercurrent-phase relation $I_1(\phi, T)$. Solid lines: $I_1(\phi, T)$ calculated for several temperatures at fixed $\epsilon = 0.05$. Dashed lines: $I_1(\phi, T = 0)$, calculated for several values of ϵ . The dotted line shows the $I_1(\phi, T = 0)$ relation which can be obtained from elementary considerations (see text).

in Refs. [3–5]). This supercurrent-phase relation can also be derived in a simple way: the jump in $I_1(\phi, T)$ at $\phi = \pi$ can be explained with Fig. 6.2, which shows that a negative state shifts below $E = 0$, and stops contributing to $I_1(\phi, T)$, whereas a positive state shifts above $E = 0$, and starts contributing to $I_1(\phi, T)$. This results in a jump [28] $\Delta I = 2ev_F/L^*$.

An increase in ϵ reduces $I_1(\phi, T)$, and the jump at $\phi = \pi$ is smoothed, until an almost sinusoidal relation is obtained at the maximum coupling $\epsilon = \frac{1}{2}$. Figure 6.3 shows that raising T has similar effects. The typical T at which $I_1(\phi, T)$ is suppressed substantially is of the order of the energy spacing between the positive and negative states. With our parameters, this corresponds to about $kT \approx 0.03\Delta \approx 0.1$ K.

We now show that $I_1(\phi, T = 0)$ can be modified by applying a voltage V between the semiconductor reservoir and the superconductors ($eV = \mu_3 - \mu_1, \mu_1 = \mu_2$). As a result electron excitations (for $V < 0$) or hole excitations (for $V > 0$) are emitted with energies $0 < E < |eV|$. The

supercurrent can now be written as [29]

$$I_1(\phi, eV) = I_1(\phi, T = 0) + \frac{1}{2} \int_0^{|eV|} J_1^{ex}(\phi, E) dE. \quad (6.10)$$

The results are shown in Fig. 6.4. The supercurrent-phase relation de-

Figure 6.4: *Supercurrent-phase relation* ($\epsilon = 0.05, T = 0$) for several values of applied voltage V between semiconductor reservoir and superconductors. The dotted lines show the extrapolated results for $\epsilon \rightarrow 0$. (a) $|eV| = n\Delta E$. (b) $|eV| = (n + \frac{1}{4})\Delta E$. (c) $|eV| = (n + \frac{1}{2})\Delta E$.

pendes periodically on the applied voltage with the period given by $\Delta E/e$, with $\Delta E = \hbar v_F/2L^*$, the spacing between consecutive positive (or negative) states. The explanation for this effect is that for energies $E < \Delta_0$, the particles from the semiconductor channel are completely reflected at the S - SM interfaces. Therefore the (non-equilibrium) occupation of the states in the channel with $E < \Delta_0$ is completely determined by μ_3 , even in the limit $\epsilon \rightarrow 0$. In real systems some inelastic scattering will always be present between the channel and the superconductors. This will tend to equilibrate the states to the electrochemical potential of the superconductors. In real systems the effect will therefore disappear in the limit $\epsilon \rightarrow 0$. The value of ϵ at which the effect disappears will be a measure of the strength of the inelastic scattering.

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- [26] L_1 and L_2 have been chosen slightly different to avoid possible effects related to the symmetry of the system.
- [27] Due to the partial reflection at the S - SM interfaces, there is a small contribution to $J_1^{ex}(\phi, E)$ due to excitations from the semiconductor reservoir with energies slightly larger than Δ_0 . This contribution vanishes when $\epsilon \rightarrow 0$.
- [28] With $V = \hbar/2e(d\phi/dt)$ one can write the supercurrent due to a fully occupied positive or negative state (which contains two electrons) as $I = 4e/\hbar(dE/d\phi) = 2ev_F/L^*$. Equation (6.8) then yields a jump $\Delta I = 2ev_F/L^*$.

- [29] When $V \neq 0$, a normal current flows between the semiconductor reservoir and the superconductors. This current vanishes for $\epsilon \rightarrow 0$.

Chapter 7

Fabrication of diffused superconducting contacts

7.1 Introduction

Almost 120 years ago Braun for the first time described the rectifying properties of metal-semiconductor contacts [1]. Physical understanding of this phenomenon, however, came only much later. In 1938 Schottky finally explained the rectification by a space charge layer on the semiconductor side of the junction which is depleted of mobile majority carriers [2, 3]. This important contribution has lead to the introduction of the term “Schottky barrier” in the physics of metal-semiconductor interfaces. The underlying physical mechanisms are, however, still under discussion and even today there is lack of a comprehensive physical theory of metal-semiconductor contacts.

In this chapter theoretical ideas will be described how to overcome the Schottky barrier in order to make ohmic contacts to a semiconductor and in particular a two-dimensional electron gas (2DEG). After that the fabrication of superconducting contacts and their characterization by means of special test structures will be discussed. Finally the measurement results and the implications for the modelling of our superconducting contacts will be treated.

7.2 Metal-semiconductor interfaces

If a metal is brought into contact with a (n -type) semiconductor, electrons will flow from the semiconductor to the metal (because of the

difference in work function $X_m - X_s$) until both Fermi energies line up. Due to this electron transfer electric fields are induced, which result in band bending. These fields also attract electrons back to the semiconductor, thus establishing the thermodynamic equilibrium. The negative charge in the metal is practically localized at the surface atomic layer (due to a very large free electron density). On the contrary, the positive charge density in the semiconductor is limited by the concentration of ionized donors and the space charge region extends into the semiconductor over a distance W_d . The band bending leads to a Schottky barrier in the semiconductor. Figure 7.1 shows the energy band diagram of a metal in contact with a n -type semiconductor (E_c is the conduction

Figure 7.1: Schematic band diagram of a metal-semiconductor interface

band edge and E_F the Fermi level).

In this simple model the height of the barrier ϕ_B can be calculated by the Schottky-Mott rule [4, 5]:

$$\phi_B = X_m - \chi_s, \quad (7.1)$$

where X_m is the work function of the metal and χ_s is the electron affinity of the semiconductor.

In practice, however, Schottky barrier heights are quite different from the energies given by this formula. Moreover, for e.g. GaAs the barrier height is not simply correlated with the work function of the metal.

Already in 1947 Bardeen [6] developed the basic concepts to resolve this discrepancy. He considered electronic states that are localized at the interface and exhibiting levels (either discretely or continuously distributed) within the energy gap. These interface states may become charge, which will influence the band bending. It is important to note that these states were postulated to be typical for the semiconductor and thus independent on the kind of metal.

Since Bardeen's publication their physical nature has remained under discussion [7]. The first conceptual model on the physical origin of the interface states was offered by Heine [8]. From the matching of wave functions at the metal-semiconductor interface he concludes [9]: *"Thus for energies in the semiconductor band gap the volume states of the metal all have tails in the semiconductor; there are actually no surface states."* He showed that the exponentially decaying tails of the metal wave functions are derived from the continuum of the virtual gap states (ViGS) of the semiconductor band structure. These ViGS are solutions of Schrödinger's equation for complex wave vectors at the centre and the boundary of the first Brillouin zone. Later these interface states were called metal-induced gap states (MIGS).

They are characterized by a neutral level ϕ_0 (measured from the valence band edge) such that the states below ϕ_0 are neutral when filled with electrons, and states above ϕ_0 are neutral when empty (see figure 7.2). In other words: at ϕ_0 the character of the states changes from predominantly donor-like to acceptor-like. It can be shown [10] that the barrier height in this model is given by

$$\phi_B = \gamma (X_m - \chi_s) + (1 - \gamma) (E_g - \phi_0) - \phi_I, \quad (7.2)$$

with

$$\gamma = \frac{\varepsilon_i}{\varepsilon_i + eN_s\delta}. \quad (7.3)$$

Here E_g is the energy gap, $\varepsilon_{i,s}$ is the permittivity of the interfacial layer, respectively the semiconductor, δ the thickness of the interfacial layer and N_s the density of interface states (per energy interval). The last

Figure 7.2: Schematic band diagram of a metal-semiconductor interface with an interfacial layer; Δ_0 is the potential across the interfacial layer

term in (7.2) ϕ_I is the image force barrier lowering, which is small in most cases.

This relation still gives a linear dependence of the barrier height as a function of the metal work function. However, the presence of interface states drastically reduces the slope parameter $\gamma = \frac{d\phi_B}{dX_m}$, which is unity in the Schottky-Mott approach. When $\gamma \rightarrow 0$, then $\phi_B \approx E_g - \phi_0$ and the Fermi level at the interface is “pinned”. For GaAs a value $\gamma \approx 0.07$ has been found, indicating a weak dependence on X_m [11]. Indeed, data for a large number of metals indicate that the intrinsic ϕ_B range for n -type GaAs is $\approx 0.6 - 0.95$ eV, independent of the contact work function, electronegativity and metal-GaAs chemical reactivity [12].

The width of the depletion region can be fairly well approximated by

$$W_d = \sqrt{\frac{2\epsilon_s(V_{bi} - V)}{eN_d}}, \quad (7.4)$$

where V is the applied voltage, N_d is the net donor concentration and the built-in voltage is given by $V_{bi} = \phi_B - E_c + E_F$.

7.3 Electronic transport through and over a Schottky barrier

For the calculation of the current across a metal-semiconductor interface two quantities have to be known as a function of energy: the supply function $N(E)$ and the transmission probability $P(E)$. The current density follows by calculating the solution of

$$j = \int_0^\infty P(E)N(E) dE. \quad (7.5)$$

The supply function is a measure of the number of incident electrons per second per unit area having energy associated with motion normal to the junction within an interval dE and it is given by

$$N(E) dE = \frac{A^*T}{k_B} \left[\ln \left(1 + e^{\frac{-(E-E_F)}{k_B T}} \right) - \ln \left(1 + e^{\frac{-(E-E_F+eV)}{k_B T}} \right) \right] dE, \quad (7.6)$$

where V is the applied bias and the Richardson constant is

$$A^* = \frac{4\pi^2 e k_B^2 m^*}{h^3}. \quad (7.7)$$

For non-degenerate semiconductors and under Maxwell-Boltzmann statistics this can be simplified to [13]

$$N(E) dE = \frac{A^*T}{k_B} e^{\frac{-(E-E_F)}{k_B T}} \left(1 - e^{\frac{-eV}{k_B T}} \right) dE. \quad (7.8)$$

For a uniformly doped semiconductor and simple parabolic bending the transmission probability can be described by [14]

$$P(E) = e^{\frac{-E_b y(\alpha)}{E_{00}}}, \quad (7.9)$$

where $\alpha = E/E_b$, $E_b = e[\phi_B - V - (E_c - E_F)]$,

$$y(\alpha) = \sqrt{1 - \alpha} - \ln \left(\frac{1 + \sqrt{1 - \alpha}}{\sqrt{\alpha}} \right) \quad (7.10)$$

and

$$E_{00} = \frac{eh}{4\pi} \sqrt{\frac{N_D}{m_r \epsilon_s}}, \quad (7.11)$$

with m_r the tunneling effective mass measured in units of the free electron mass and N_D the donor concentration. This can be approximated for small energy by [15]

$$P(E) \approx e^{-\frac{E_b - E}{E_{00}}} \approx e^{-\frac{E_b}{E_{00}}}. \quad (7.12)$$

For GaAs $E_{00} \approx 2.0 \times 10^{-14} \sqrt{N_D}$ eV.

The expression for the current (7.5) can be divided into two components. One is due to thermionic emission (TE) of electrons over the barrier; the other is due to tunneling of electrons through the barrier. Tunneling taking place at an energy near E_F is called field emission (FE); thermally assisted tunneling, above E_F , is called thermionic field emission (TFE). The thermionic component is approximately

$$j_{TE} = A^* T^2 e^{-\frac{e\phi_B}{k_B T}} \left(e^{\frac{eV}{k_B T}} - 1 \right) \equiv j_{sat} \left(e^{\frac{eV}{k_B T}} - 1 \right). \quad (7.13)$$

The contribution of (thermionic) field emission can not easily be calculated analytically. Series solutions can be obtained under specific conditions. An important quantity turns out to be E_{00} , which is mainly determined by the band bending at the interface. When $k_B T \sim E_{00}$ thermionic field emission is dominant and the following expression is approximately valid:

$$j_{TFE} = j_0 e^{\frac{eV}{E_0}} \text{ with } E_0 = E_{00} \coth \frac{E_{00}}{k_B T}. \quad (7.14)$$

The quantity j_0 is given by a complicated expression. In GaAs thermionic field emission occurs roughly for $N_D > 10^{23} \text{ m}^{-3}$ at 300 K and for $N_D > 10^{22} \text{ m}^{-3}$ at 77 K.

At very high doping levels ($k_B T \ll E_{00}$) the width of the depletion region becomes so narrow that field emission is dominant. The current-voltage characteristics in this regime are described by

$$j_{FE} = j_s e^{\frac{eV}{E_{00}}}. \quad (7.15)$$

By means of the equations 7.13, 7.14 and 7.15 the specific contact resistance defined as

$$\rho_c = \left(\frac{\partial j}{\partial V} \bigg|_{V=0} \right)^{-1} \quad (7.16)$$

can be calculated. For the regime of thermionic emission is found:

$$\rho_{c,TE} = \frac{k_B T}{q A^* T^2} e^{\frac{\phi_B}{k_B T}} \quad (7.17)$$

and for field emission:

$$\rho_{c,FE} = A e^{\frac{\phi_B}{E_{00}}}, \quad (7.18)$$

where A is a function of E_{00} , ϕ_B , $E_c - E_F$ and T .

It is important to note that $\rho_{c,TE}$ is only dependent on the barrier height ϕ_B and temperature, while $\rho_{c,FE}$ is also dependent on E_{00} . From this analysis it can be concluded that the design of low resistance contacts requires low barrier height ϕ_B and high doping.

The problem in GaAs is that ϕ_B cannot be made small by the proper choice of metal, because of a high interface-state density. The alternative then is to make the barrier as thin as possible, so that the current transport through the barrier is enhanced by tunneling (the field emission regime). However, for n -type GaAs, there are at least three types of upper limits on donor concentration [16]: (1) the most commonly used donor species (group IV elements) are amphoteric, (2) DX centres create resonant states in the conduction band which limit free electron concentration and (3) dopant precipitation and phase formation may occur due to solid solubility limitations. Because of these difficulties and the usually more complicated metallurgy the development of ohmic contacts to GaAs has been delayed relative to that for contacts to Si.

7.4 Ohmic contacts to GaAs

For practical purposes an ohmic contact can be defined as a source of carriers with an internal resistance R_c which is totally negligible compared with the semiconductor resistance or as a source of carriers with a non-negligible internal resistance R_c , but one which obeys Ohm's law for current densities of interest [17].

As already mentioned in the previous section, for GaAs it is difficult to reduce the barrier ϕ_B in order to obtain a low contact resistance. Nevertheless, contacts have been made by growing an intermediate layer between the metal and GaAs. If the interlayer is compositionally graded and the barrier between the metal and the interlayer is negligible, a contact with no barrier to current flow can be formed. This has been successfully accomplished with a graded $n\text{-In}_x\text{Ga}_{1-x}\text{As}$ interlayer [18].

The grading from InAs to GaAs effectively compensates the electric field which would exist at an abrupt InAs/GaAs interface. This model may offer an explanation of why pure In has always proved to be a good contact for n -type GaAs, even though it is not a donor.

In practice, a low resistance ohmic contact is usually obtained by increasing the doping near the metal-semiconductor interface. The band diagram of such a metal/ n^+ -GaAs/ n -GaAs contact is shown in figure 7.3. Often the the highly doped region near the surface is created

Figure 7.3: *Band diagram of a metal/ n^+ -GaAs/ n -GaAs contact*

by ion implantation or diffusion. The doping concentration is limited by the impurity solubility and may reach $5 \times 10^{25} \text{ m}^{-3}$ for n -type GaAs. The observed low values of contact resistivity can however not be explained satisfactorily by the tunneling model. For the most popular metallization, composed of Au, Ge and Ni, Braslau [19] shows that typical values of ρ_c are much lower for a given N_D than would have been predicted. Moreover, it appears that $\rho_c \propto N_D^{-1}$, rather than $\log \rho_c \propto \sqrt{N_D}^{-1}$. To explain this a model has been formulated to take account of highly doped spikes, protruding into the GaAs [19].

As an example, the most widely used contact system AuGeNi will shortly be described now. Although it is also the most investigated system, it is still not completely understood. This metallization was invented in 1967 [20]. Earlier relative low contact resistances were obtained by annealing AuGe contacts with eutectic composition (Au-13wt%Ge) at $T \approx 360^\circ\text{C}$. Significant improvement of the binary AuGe contacts was made by adding Ni (first from the contamination of the tungsten evap-

erator source...) to the AuGe: the surface roughness was reduced. By depositing a thin first layer of Ni the adhesion was improved, together with reproducibility and device reliability. Lowest contact resistances are obtained at annealing temperatures above 440°C

XRD and TEM examinations¹ have shown that NiAs(Ge), β -AuGa and other compounds are formed at the interface of the contacts. Direct contact of the NiAs(Ge) phase with the GaAs interface is important for low contact resistances. The most widely accepted model [16] is that this low contact resistance is due to heavy Ge doping at the interface. During contact annealing Ga atoms diffuse out from the GaAs substrate to the contact material (forming β -AuGa phases), creating Ga vacancies, while Ge diffuses from the NiAs(Ge) phase to the Ga vacancy sites, providing donors close to the metal-GaAs interface at a level higher than $5 \times 10^{24} \text{ m}^{-3}$. This narrows the depletion region and allows tunneling of electrons through the barrier.

The β -AuGa compounds are essential for reduction of the contact resistance, since the for doping necessary Ga vacancies are created by forming β -AuGa at the interface. However, these phases have a negative effect on the thermal stability after contact formation.

The surface roughness can be reduced partially by covering the contact materials by top cap layers.

7.5 Superconducting contacts

Our goal was to make a superconducting contact to the two-dimensional electron gas (2DEG) in GaAs/AlGaAs heterostructures. To achieve that the superconductivity extends into the 2DEG, after contact annealing either all contact layers should be superconducting, or eventual normal layers should be so thin that the 2DEG is phasecoherently connected to the superconducting regions. The fact that the 2DEG is situated typically 100 nm below the surface of the heterostructure makes this a severe requirement.

In the literature no research on this specific subject could be found. Only three experiments had been described that are in some way related.

Superconducting shallow AuGe/Nb contacts have been made to n -GaAs by Gurvitch *et al.* [21]. The best results were obtained for 15 nm AuGe covered with 200 nm Nb, annealed in reducing atmosphere at

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¹XRD is X-Ray Diffraction, TEM is Transmission Electron Microscopy

390–420°C for 1–5 s. In this system clearly Nb is the superconductor; the ohmic contact has to be formed by Ge doping. Nb apparently also serves as a hard refractory overlayer which holds the molten eutectic in place and prevents it for “balling up”. Although indications for a superconducting transition has been observed at $T \approx 8$ K, it is by no means clear from the paper, whether, and if so, how far superconductivity is induced into the semiconductor.

In 1987 Ivanov *et al.* [22] reported that superconducting Sn source and drain electrodes had been weakly connected via thinned ohmic contacts and an accumulation layer induced at an $\text{Al}_{0.5}\text{Ga}_{0.5}\text{As}/n\text{-GaAs}$ insulator/semiconductor interface. The ohmic contacts consisted of 100 nm In which was alloyed at 450°C during 30 s. It was necessary to etch away “most of the ohmic contact region” by Ar ion bombardment before the superconducting Sn electrodes were deposited (100 nm). Since no doping has been introduced, the operation of this ohmic contact must be based on the band bending caused by alloying In (compare also the previous section). Probably this alloy was not superconducting itself and therefore it had to be etched away mainly to establish phase coherence between the accumulation layer and the superconductor (Sn).

In 1989 Ivanov described superconducting properties of alloyed ohmic In contacts to $n\text{-GaAs}$ [23]. During alloying the 200 nm In layer is covered by a 200 nm SiO_2 cap layer. This is necessary to reduce surface roughness and to obtain low resistances and reproducible electric characteristics. Analogously low-resistance Sn contacts have been made; however, their superconducting properties have not been investigated separately.

In principle it is recommendable to use only superconductors for the contacts. Of all superconducting materials only the elements In and Sn are known to form low-resistance ohmic contacts to GaAs. The In contacts are based on band bending by alloying which seems less controllable in practice (especially in heterostructures). Besides that, In is an inconvenient material to use in the fabrication of nanostructures because of its low melting temperature of 156°C, which is, for example, below the usual baking temperature of PMMA resist. Therefore we preferred to use Sn.

In the past Sn contacts were made by annealing macroscopic Sn pieces on top of the GaAs/AlGaAs heterostructure. For submicron devices it is however necessary to use patterned contacts fabricated from evaporated thin films. This is by no means trivial, since Sn films im-

#	Ref.		A		B		C		D		E	
4	Au	50	Au	50	Au	50	Au	50	Sn	80		
3	Ni	55	Nb	55	Nb	55	Ti	55	Ti	55		
2	AuGe	170	AuGe	55	Sn	170	Sn	170	Sn	170	Sn	200
1	Ni	12	Nb	12	Nb	12	Ti	12	Ti	12	Ti	12

Table 7.1: *The layers of the different contact metallizations with thickness in nm*

mediately “ball up” and form droplets when heated, instead of diffusing into the semiconductor. In 1990 we found that this could be prevented by placing a 12 nm Ti layer between the GaAs and the Sn. This intermediate layer provides a better wetting of the semiconductor surface and acts as a “sieve” for the tin. Titanium is a very nice material for this purpose, since it has a high melting temperature (1677°C), it is a superconductor ($T_c = 0.4$ K), its atomic radius is approximately equal to that of Ga and it forms a relatively modest Schottky barrier on GaAs ($\phi_B = 0.83$ eV [12]). In contrast with the SiO₂ cap layer of Ref. [23] this layer doesn’t have to be removed after the annealing.

The first contacts made by this process (12 nm Ti/200 nm Sn annealed at 405°C during 5 s in N₂ atmosphere) showed clear evidence of superconductivity (see the following chapter). Although our primary goal, i.e. to investigate whether it is possible to induce superconductivity into the 2DEG, was achieved, it is useful trying to optimize and characterize these contacts.

7.6 Contact research

Considering the function of the different layers in the AuGeNi contact system and the successful superconducting contacting with Ti/Sn, a plan for further investigation and optimization of diffused superconducting contacts has been made. For this purpose 6 different contact metallizations have been annealed during 4 different times. The layers of the contact schemes are shown in table 7.1.

The Ni/AuGe/Ni/Au metallization is the commonly used normal ohmic contact in our group; in this research it serves as a reference contact for the superconducting contacts. Material system **A** is chosen

as a superconducting analog of the normal reference. Instead of Ni (which is magnetic and thus has to be avoided) the superconductor Nb is used which has also a high melting temperature (2497°C). Just like Ni it should improve adhesion to GaAs and surface roughness. The eutectic AuGe provides the donors for doping the semiconductor and creates Ga vacancies. The top layer Au prevents the contact from oxidation during annealing and provides a good surface to attach bond wires to. As discussed in ref. [21] the thickness of the AuGe has to be kept small, since otherwise superconductivity cannot be induced throughout this layer by Nb.

In contact **B** the AuGe is replaced by Sn. Now all essential layers of the contact are superconductors itself. From experience with the Ti/Sn contacts we know that the Sn layer should be not too thin, therefore it is made 170 nm thick (comparable to the thickness in the Ti/Sn contacts)²; indeed, it is no problem if part of the Sn layer is left after annealing, since it is a superconductor.

The metallization schemes **C** and **D** are variations of the Ti/Sn contact. On top a Ti layer and, respectively, an Au or Sn layer are evaporated. It was expected that this would improve surface roughness and facilitate bonding. Besides that, the second Sn layer could maybe serve as an extra source of dopant. Finally, of course the original Ti/Sn metallization is added as contact **E**.

As can be seen in table 7.1 in all contacts thicknesses of corresponding layers are kept the same as much as possible.

For each of the six metallizations five samples have been made on one and the same GaAs/Al_{0.33}Ga_{0.67}As wafer (W222). In this material the 2DEG, with electron mobility $\mu \approx 110 \text{ m}^2/\text{Vs}$ and electron density $N \approx 3 \times 10^{15} \text{ m}^{-2}$ ($T = 5 \text{ K}$), is situated $\approx 95 \text{ nm}$ below the surface.

First the surface of the heterostructure is cleaned: 2 minutes in warm acetone of 60°C followed by ultrasonic agitation during 3 minutes, 1 minute at 60°C in trichloroethylene:acetone:isopropanol (1:1:1), rinsing in acetone. After this the heterostructure is dried by spinning.

Now the (positive) resist layers of PMMA (polymethylmetacrylate) are applied by spinning. The specifications are shown in table 7.2.

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²It has to be noted that the maximum total thickness of the metallization is limited by the nanolithography process. It is not easy to make the e-beam resist (PMMA) layer very thick and still homogeneous; moreover with increasing thickness resolution is lost. In this research the PMMA double-layer was $\approx 600 \text{ nm}$ thick. Although for lift-off the thickness of the evaporated metal film is usually kept below one third of that of the (lower) resist layer, we found that up to 320 nm is possible in our system.

layer #	molecular mass	solvent	baking time at 175°C	resulting thickness
1	180K	chlorobenzene	1 hour	≈450 nm
2	350K	ortho-xylene	1 hour	≈150 nm

Table 7.2: Specifications of the used PMMA layers

By means of a 50 keV electron beam so-called TLM-structures³ (see also next section) are defined in this layer with a dose of $\approx 475 \mu\text{C}/\text{cm}^{-2}$. The TLM-structure consists of 6 $200 \mu\text{m} \times 200 \mu\text{m}$ squares in a line on mutual distances of respectively $800 \mu\text{m}$, $400 \mu\text{m}$, $200 \mu\text{m}$, $100 \mu\text{m}$ and $60 \mu\text{m}$ (see figure 7.4). The patterns are developed during 1 minute in

Figure 7.4: Layout of the used TLM-structure (on scale)

isopropanol:MIBK⁴ (3:1) and rinsed in pure isopropanol.

In this point the samples are ready for metal evaporation. Immediately before each sample is placed in the vacuum chamber of the evaporation equipment, it is etched in $\text{NH}_4\text{OH}(25\%):\text{H}_2\text{O}$ (1:25) during 1 minute to remove oxides. After the etch the sample is successively rinsed in demineralized water, blown dry and mounted in the evaporation chamber. When the pressure has decreased below 10^{-5} Pa, the metal films are deposited perpendicularly by means of e-gun evaporation. Since in particular Ti and Nb are known to “getter”, all sources are cleaned by pre-heating before the shutter is opened. Moreover, it is important to evaporate Ti and Nb with a high rate ($\gtrsim 1 \text{ nm}$) in order to obtain non-oxidized and superconducting films. The evaporation rates for the different metals are shown in table 7.3. After the metallization “lift off” takes place in boiling acetone (with ultrasonic agitation).

To obtain a well-specified geometry of the sample a mesa is defined by optical lithography. Using the positive photo-resist AZ S18.13 and

³TLM means Transmission-Line Model

⁴MIBK is methylisobutylketone

metal	rate (nm/s)
Au	0.2
AuGe	0.2
Ni	0.5
Sn	0.5
Ti	$\gtrsim 1$
Nb	$\gtrsim 2$

Table 7.3: Overview of the evaporation rates for the used metals

UV-exposure a bar ($200\text{ }\mu\text{m} \times 2.76\text{ mm}$) is patterned over the metallic squares. The 2DEG outside this bar is removed by etching in $50\text{ ml H}_2\text{SO}_4(98\%):\text{H}_2\text{O}(1:5) + 2\text{ ml H}_2\text{O}_2(30\%)$ during 20 s. This yields an approximately 170 nm high mesa⁵. Now the samples are cleaved and are ready for annealing.

Besides studying the variation of contact materials, we also wanted to get information about the optimal annealing time. In order to do that reproducible heating during short intervals ($\sim 10\text{ s}$) is necessary. Annealing on a hot plate (which is normally used) is not reproducible enough for this purpose. “Rapid thermal annealers” are very reproducible, but the cooling down time is usually a few minutes. This would increase the effective annealing time beyond the range which is relevant for our contacts. The best compromise was found to be a strip-heater. The annealing process can be programmed very reproducibly by means of electronic control systems. In this way a heating curve as shown in figure 7.5 can be obtained. Only the annealing time t_{ann} is varied, while all other parameters are kept constant. Suitable annealing temperatures are already known ($T_{\text{ann}} = 450^\circ\text{C}$ for contacts containing AuGe and $T_{\text{ann}} = 400^\circ\text{C}$ for contacts containing Sn), and variation of it would give rise to complicated differences in metallurgical and chemical processes.

We can define the effective annealing time t_{eff} as the period during which the temperature is above the melting temperature T_m of AuGe

⁵On some samples the edges of the metal squares were slightly damaged by this etch, which can be avoided by choosing a different order of processing steps (as is done for all samples described in other chapters). However, in this case we wanted to process all samples of a certain kind as long as possible together to enable a good comparison.

Figure 7.5: *Programmed heating curve of the strip-heater (schematically); T_{ann} is the desired annealing temperature, T_m the melting temperature, t_{ann} is the programmed time at the plateau and t_{eff} the resulting effective annealing time*

(370°C) respectively Sn (231°C). As can be seen in the figure the effective annealing time is a small, fixed amount of time longer than the programmed time t_{ann} ; this extra time is approximately 3 s for AuGe-based metallizations and 4 s for Sn-based contacts.

For all different contacts a sample is annealed during $t_{\text{ann}} = 0, 6, 18$ and 300 s. One sample of each kind is kept as a reference. The annealing takes place in an atmosphere of forming gas (N_2/H_2); care is taken to place the sample in the centre of the strip to avoid temperature gradients. The temperature is measured by a thermocouple directly under the sample. Since the samples are small ($\sim 3 \times 5$ mm) the heat capacity will be small.

Finally two Al wires are attached to each contact by means of ultrasonic bonding to make four-terminal measurements possible.

7.7 Characterization of ohmic contacts

It can be very complicated to get information about contact properties in a semiconductor system. Therefore special methods have been developed for this purpose; one of these is the analysis by the transmission-line

model (TLM).

The TLM structure, used for planar contact-resistance measurements is a series of identical ohmic contacts, of width w and length l_c , spaced at varying intervals $l_1, l_2, l_3, \dots$ (like the pattern in figure 7.4). The resistance measured between two neighbouring contacts is given by:

$$R = 2R_c + r_s \frac{l_i}{w}, \quad (7.19)$$

where R_c is the contact resistance and r_s the sheet resistance of the semiconductor material between the contacts. By plotting the measured resistance at different l_i the slope of the resulting line will be r_s/w and the intercept with the R -axis $2R_c$. Thus the TLM-structure provides useful information which enables a comparison between our different contact types and with the reference Ni/AuGe/Ni/Au contact. This will be discussed in section 7.8.

If information about the specific contact resistance ρ_c is wanted, a much more detailed analysis is needed and a suitable model has to be found for the contact region, which can be very complex in reality [17]. In figure 7.6 two (extreme) possibilities for a model of an annealed ohmic contact to a bulk semiconductor are shown. In the left picture

Figure 7.6: Schematic diagrams of possible scenarios for annealed ohmic contacts: (a) the metal diffuses through the active semiconductor layer and retains its metallic character; (b) one or more elements of the metal diffuse through the layer, but merely act as semiconductor dopants

the diffused (or alloyed) metal goes through the entire active region of thickness d . Then the current flows uniformly out of a contact area wd ,

and $\rho_c = R_c w d$. However in practice, the other possibility is often more plausible: the metal remains nearly planar, but some elements (e.g. Ge or Sn) diffuse partially into the semiconductor. Now the current won't flow uniformly out of the entire metal area, but will crowd toward the edge nearest to the other contact. This can be modeled in terms of a transmission line [24, 25] (see figure 7.7).

Figure 7.7: Current flow in a thin-layer semiconductor device in terms of the transmission-line model

Consider the current flow as in the figure. A current $I(0)$ which enters the contact region at $x = 0$, flows horizontally under the metal with various portions being lost in vertically direction through the infinitesimal vertical resistors $\rho_c/w \, dx$. This process obeys the following equations:

$$dV(x) = -\frac{R_s}{w} I(x) dx \quad (7.20)$$

$$dI(x) = \frac{w}{\rho_c} (V_c - V(x)) dx, \quad (7.21)$$

where R_s is the sheet resistance of the material *under* the contact and V_c is the potential in the metal layer of the contact. Solving these equations by means of Laplace transformation gives an expression for the contact resistance:

$$R_c \equiv \frac{V(0) - V_c}{I(0)} = \frac{R_s l_T}{w} \coth \frac{l_c}{l_T}, \quad (7.22)$$

where the transfer length $l_T \equiv \sqrt{\frac{\rho_c}{R_s}}$. This is the typical distance from the edge over which current enters the contact. Thus the effective area for a planar contact is circa wl_T instead of wl_c .

For electrically long contacts (i.e. $l_c/l_T \gtrsim 2$, which in practice holds for nearly all contacts) is

$$\rho_c \approx \frac{w^2 R_c^2}{R_s}. \quad (7.23)$$

and the transfer length is

$$l_T = \frac{w R_c}{R_s}. \quad (7.24)$$

Experimentally, however, the value of the sheet resistance R_s of the semiconductor *under* the contact is usually not known. Often R_s is assumed to be equal to $R_{\square}$, the sheet resistance of the semiconductor *between* the contacts, which is certainly not always correct.

If the active semiconductor region under the TLM-pattern is not homogeneous in depth, the described “one-layer TLM” analysis has to be adapted. Especially for application to heterostructures a “two-layer TLM” has been developed [17] (see figure 7.8). In this model besides

Figure 7.8: Current flow in a two-layer semiconductor device in terms of the transmission-line model

conduction through the 2DEG (with sheet resistance $R_{\square 2}$), parallel current flow through the doped AlGaAs layer (with sheet resistance $R_{\square 1}$) is taken into account. The interface between the metal and the AlGaAs

and the AlGaAs/GaAs interface are modeled as barriers with specific contact resistivities ρ_{c1} and ρ_{c2} respectively. For our samples, however, we expect that the current flow through the AlGaAs layer is negligible compared with that through the 2DEG, since $R_{\square 1} \gg R_{\square 2}$. In practice the small series resistance of the doped three-dimensional semiconductor region between the two barriers under the contact can be ignored. Therefore we can use the one-layer transmission-line model to analyze our devices, with ρ_c containing both ρ_{c1} and ρ_{c2} in figure 7.8.

7.8 Results of the electrical characterization

First the $I - V$ -characteristics of all neighbouring contact pairs on all 24 samples have been measured at room temperature, using a standard voltage-biased DC technique ($-5 \text{ V} < V < 5 \text{ V}$). By means of these curves the resistance around zero-bias can be determined and it can be checked whether the resistance behaves linearly.

Since the resistance of the metal region between both leads on a contact ($\sim 1\Omega$) is much smaller than the resistance of the leads ($\sim 10^2 \Omega$) and the 2DEG resistance ($\sim 10^3 \Omega$) the four-terminal resistance can be obtained from four two-terminal measurements. For two contacts (1 and 2), each with two leads (A and B) we find the following expressions for the possible resistance measurements:

$$R_{1A-1B} = R_{1A} + R_{1B} \quad (7.25)$$

$$R_{2A-2B} = R_{2A} + R_{2B} \quad (7.26)$$

$$R_{1A-2A} = R_{1A} + R_{2A} + 2R_c + \frac{l_i}{w} \cdot R_{\square} \quad (7.27)$$

$$R_{1B-2B} = R_{1B} + R_{2B} + 2R_c + \frac{l_i}{w} \cdot R_{\square} \quad (7.28)$$

and so

$$\begin{aligned} R &= 2R_c + \frac{l_i}{w} \cdot R_{\square} \\ &= \frac{1}{2}(R_{1A-2A} + R_{1B-2B} - R_{1A-1B} - R_{2A-2B}). \end{aligned} \quad (7.29)$$

Here $R_{\square}$ denotes the specific resistance of the 2DEG.

If this resistance R is plotted against the contact separations l_i , R_c and $R_{\square}$ can be determined as described in section 7.7. These plots are shown in figure 7.9 for the sample with the lowest resistance in each

Figure 7.9: The resistances of the TLM-structure against the distances l_i for one sample of each contact type ($T = T_{\text{room}}$)

series. The measured values of the averaged resistances are presented in table 7.4.

From the slope of the plots a value for $R_{\square}$ is found around $2 \text{ k}\Omega$, which is in reasonable agreement with the specified value $(n_e e \mu)^{-1} = 1.6 \text{ k}\Omega$. The slope is practically the same for all samples, which indicates that the 2DEG between the contacts is not affected by the contact processes.

A very rough estimation for the contact resistivity ρ_c can be made by assuming that $R_s = R_{\square}$. The thus obtained values are presented in table 7.5; to account for the l -independent resistance between the two-dimensional and the three-dimensional region of the contact, as an estimation for an under limit a Sharvin resistance $R_{\text{Sh}} \equiv \frac{h}{2e^2} \cdot \frac{\lambda_F}{2w} \approx 1.5 \Omega$ has been subtracted, serving as an estimation for an under limit. In this way only an upper limit for ρ_c is obtained; in particular, it is expected that R_s is (much) larger than $R_{\square}$. However, the data of the

	$t_{\text{ann}}=0$ s	$t_{\text{ann}}=6$ s	$t_{\text{ann}}=18$ s	$t_{\text{ann}}=300$ s
Ni/AuGe/Ni/Au	24	22	17	28
Nb/AuGe/Nb/Au	55	52	27	56
Nb/Sn/Nb/Au	$2 \cdot 10^3$	430	770	270
Ti/Sn/Ti/Au	$25 \cdot 10^3$	$2.7 \cdot 10^3$	$1.3 \cdot 10^3$	$1.2 \cdot 10^3$
Ti/Sn/Ti/Sn	$1.5 \cdot 10^3$	$3.7 \cdot 10^3$	510	$2.2 \cdot 10^3$
Ti/Sn	140	120	215	100

Table 7.4: Observed contact resistances (in Ω) at room temperature

	$t_{\text{ann}}=0$ s	$t_{\text{ann}}=6$ s	$t_{\text{ann}}=18$ s	$t_{\text{ann}}=300$ s
Ni/AuGe/Ni/Au	$9.8 \cdot 10^{-4}$	$8.5 \cdot 10^{-5}$	$5.1 \cdot 10^{-5}$	$1.5 \cdot 10^{-4}$
Nb/AuGe/Nb/Au	$5.9 \cdot 10^{-4}$	$5.1 \cdot 10^{-4}$	$1.3 \cdot 10^{-4}$	$6 \cdot 10^{-4}$
Nb/Sn/Nb/Au	$8.2 \cdot 10^{-1}$	$3.8 \cdot 10^{-2}$	$1.2 \cdot 10^{-1}$	$1.6 \cdot 10^{-2}$
Ti/Sn/Ti/Au	125	1.5	$3.2 \cdot 10^{-1}$	$3.0 \cdot 10^{-1}$
Ti/Sn/Ti/Sn	$4.5 \cdot 10^{-1}$	2.8	$5.3 \cdot 10^{-2}$	$9.8 \cdot 10^{-1}$
Ti/Sn	$3.9 \cdot 10^{-3}$	$2.9 \cdot 10^{-3}$	$9.3 \cdot 10^{-3}$	$2 \cdot 10^{-3}$

Table 7.5: Rough estimations for the specific contact resistances ρ_c (in Ωcm^2) at room temperature

contact Ni/AuGe/Ni/Au serve as a reference for comparison of the other contacts, which is sufficient for our purpose.

A first requirement for good ohmic contacts (normal or superconducting) is a sufficiently low (and ohmic) contact resistance at room temperature. On cooling down R_c will in principle increase, since thermionic emission is suppressed and the effect of an interface barrier will become more pronounced. This is especially important for superconducting contacts, since Andreev reflection is a two-particle process and thus the probability for this process will have a square dependence on the transmission of the barrier. Moreover, if the semiconductor below the contact is not doped sufficiently, it will “freeze out” and contact with the 2DEG is lost.

It is clear from tables 7.4 and 7.5 that the resistances of the Ti/Sn/Ti/Au and the Ti/Sn/Ti/Sn contacts are much higher and less constant

	$t_{\text{ann}}=0$ s	$t_{\text{ann}}=6$ s	$t_{\text{ann}}=18$ s	$t_{\text{ann}}=300$ s
Nb/AuGe/Nb/Au	70	60	14	10
Nb/Sn/Nb/Au	$\approx 40 \cdot 10^3$	$\approx 30 \cdot 10^3$	1500	30
Ti/Sn	20	35	925	$> 40 \cdot 10^3$

Table 7.6: Observed contact resistances R_c (in Ω) at $T = 1.2$ K

than those of the other contacts. Their $I - V$ -curves are also strongly non-linear, which shows that a significant Schottky barrier still exists at the interface. This means that these metallizations are not suitable for superconducting ohmic contacts.

To study the properties of the contacts at lower temperatures the samples with potential superconducting contacts are mounted in a helium cryostat; here the temperature can be varied from 4.2 K to 1.2 K. The contact resistances at $T = 1.2$ K are collected in table 7.6. The $dV/dI - V$ -curves of most contacts show peak structures around zero-bias. These are not included in the R_c -values; to calculate the resistances in this table the resistance at voltages above the peak is extrapolated to a value at zero-bias.

The resistances of the Nb/Sn/Nb/Au contacts are very high and not a symmetric function of voltage for all but the largest t_{ann} . It can be concluded that not yet an *ohmic* contact has been formed. The contacts which were annealed during 300 s have completely flat $dV/dI - V$ -curves at low biases. Nevertheless, it has been proved by resistance measurements through a single contact as a function of magnetic field that superconductivity still exists. However, at $T = 1.2$ K no indications could be found of Andreev reflection.

More symmetric resistance peaks around zero voltage are found in the Nb/AuGe/Nb/Au samples. The dependence of resistance on magnetic fields up to 0.2 T is however very weak, so that it is not clear whether they can be caused (completely) by superconductivity. In the sample annealed during $t_{\text{ann}} = 300$ s structures which could indicate enhanced weak localization (see chapter 10) have been found for the smallest separation of 60 μm .

Of the Ti/Sn contacts the samples of $t_{\text{ann}} = 18$ s and $t_{\text{ann}} = 300$ s have much higher and less constant contact resistances than for the other two annealing times. For $t_{\text{ann}} = 0$ s and $t_{\text{ann}} = 6$ s broad resistance peaks

are found around zero bias (see figure 7.10). These can be explained

Figure 7.10: *The differential resistance of the Ti/Sn sample with $t_{\text{ann}} = 0$ s ($l = 200$ μm) as a function of the bias current at $T = 1.2$ K*

well by the opening of the superconducting gap of Sn, which introduces Andreev reflection, in combination with a small barrier [26]. Below the critical temperature of bulk Sn ($T_c = 3.7$) the resistance starts to increase with decreasing T and a peak is formed in the $dV/dI - V$ -curves (see figure 7.11). At $T = 1.2$ K this peak extends on the voltage scale up to approximately twice the superconducting gap Δ of Sn ($2\Delta \approx 1.14$ meV). The origin of the two little dips on the slopes of the peak that can be seen in figure 7.10 is not clear, since they are only observed for $l = 200$ μm on the sample with $t_{\text{ann}} = 0$ s. The shape of the broad peak is however practically the same for the other contacts.

This peak is also suppressed at a magnetic field of ≈ 35 mT, which is in a good agreement with the critical field for bulk Sn; at the same value the two-terminal resistance of a single contact begins to increase, indicating the suppression of superconductivity.

The contact resistances of these samples are low. Moreover they seem to be stable in time: after a year (at room temperature) the resistance is practically the same. Together with the clear indications for superconductivity this makes the Ti/Sn contacts the most preferable of

Figure 7.11: *The zero-bias differential resistance of the Ti/Sn sample with $t_{\text{ann}} = 0$ s ($l = 200 \mu\text{m}$) as a function of the temperature*

the investigated metallization schemes.

7.9 Discussion and conclusions

It is interesting to note that from superconducting materials contacts to a 2DEG can be made with contact resistances that are comparable to that of the commonly used Ni/AuGe/Ni/Au metallization. In this respect Nb/AuGe/Nb/Au, but also Ti/Sn give remarkable results, both at room temperature and at 1.2 K. Above temperatures where superconductivity plays a role the resistance of these contacts behaves completely ohmic in the considered current range (up to $\sim 10^{-3}$ A).

The optimum annealing times are different per contact type (see table 7.7). Since in general no sharp increase in contact resistance is observed for $t_{\text{ann}} = 300$ s, the metallizations seem to contain abundant doping material; the diffusion into the semiconductor is not limited by a finity of the source. Only the resistance of the Ti/Sn contact with $t_{\text{ann}} = 300$ s might be improved by a thicker Sn layer (see table 7.6).

The “smart” variations Ti/Sn/Ti/Au and Ti/Sn/Ti/Sn yield seri-

contact	$t_{\text{ann}}^{\text{opt}}$ (in s)
Ni/AuGe/Ni/Au	18
Nb/AuGe/Nb/Au	300
Nb/Sn/Nb/Au	300
Ti/Sn/Ti/Au	300
Ti/Sn/Ti/Sn	18
Ti/Sn	0

Table 7.7: Recommended annealing times t_{ann} for the different contact types; $T_{\text{ann}} = 400^\circ\text{C}$ for Sn-based metallizations and $T_{\text{ann}} = 450^\circ\text{C}$ for Ge-based metallizations

ously worse instead of improved contact properties. The difference with Ti/Sn may be that part of the Sn diffuses into the overlaying metals instead of into the semiconductor. If that is true, it would be worthwhile to try Nb/Sn as an alternative for Ti/Sn, since the same effect could be important in the Nb/Sn/Nb/Au metallization.

The Nb/AuGe/Nb/Au contacts have very low resistances, but at 1.2 K no clear signatures of superconductivity have been observed. Maybe variation of the layer thicknesses (for example, a thicker upper Nb layer or an even thinner AuGe layer) can improve this.

The contact resistances of the Nb/Sn/Nb/Au contacts ($T = 1.2\text{ K}$) for lower t_{ann} are too high to be useful. The contacts annealed during 300 s are much less resistive and are at least partially superconducting. At this annealing time, however, elements can diffuse far into the semiconductor, which is unwanted for use of the contacts in nanostructures.

The most preferable contact scheme is the original Ti/Sn metallization. It combines a low contact resistance with clear superconducting effects. Moreover the optimum annealing time is short. That in these experiments a resistance peak is observed in the $dV/dI - V$ -curves at 1.2 K instead of a reduction in the resistance, is probably due to the specific properties of the GaAs/AlGaAs heterostructure (e.g. n_e) that is used. If we compare this with samples on different heterostructure material, which are described in the following chapters, we see that the exact behaviour of the $dV/dI - V$ -characteristics is dependent on the kind of used semiconductor material, although the normal contact resistance is always approximately the same. Nevertheless, all contacts made

on one kind of material behave in a comparable way, so the annealing process itself seems to be well reproducible.

Of course we would like to know more about the way in which these Ti/Sn contacts function. Because the contacts don't "freeze out", the semiconductor between metal and 2DEG has to be heavily doped. From the fact that the contact resistances at $T = 1.2$ K are in general not higher than at room temperature, we can conclude that the contacts are in the transport regime where field emission is dominant. It is difficult to develop a precise model of the contact region. Since in the superconducting state a resistance peak is observed, the transmission of the NS interface is not unity. However, since in the normal state R_c is completely ohmic and very weakly dependent on T an eventual barrier have to be extremely transmissive. This paradoxical situation is possible, since Andreev reflection has a square dependence on the transmission of a barrier, in contrast with normal electron tunneling which has a linear dependence. Because resistance dips have been observed on other heterostructures, the requirements on electron density of the 2DEG could be critical in this respect. In that case it would be expected that not a barrier at the metal/GaAs interface but one at the original AlGaAs/GaAs interface is the limiting factor. Indeed, the width of this depletion region is directly dependent on n_e .

Some researchers recently have suggested on the base of experimental indications that still a 2DEG exists under the contact region [17,27,28]. If this is true, it would per definition imply an obstruction between the 3D contact region and the 2DEG. Detailed magnetoresistance measurements and a systematic study of the influence of n_e on R_c can probably give more insight. In this respect it would be interesting to make the same contacts to a so-called inverted heterostructure, where the doped AlGaAs is not situated between the contact and the 2DEG, but on the other side of the 2DEG.

For superconducting contacts it is extremely important that the contact region is phase-coherent. Another idea to increase the transmission of the contact region is to use a shallow 2DEG. Since here the distance d from metal to 2DEG is shorter, it should be easier to achieve phase coherence down to the 2DEG. We already have made samples on such material, with $d \approx 26$ nm, (see chapter 9), but the transmission in this specific case was not as high as hoped. The reason for this is probably the different layer structure of that material. Between the surface and the 2DEG it contains two AlAs layers, which can have a negative ef-

fect. In the first place, the barrier at an AlAs/GaAs interface is $\approx 3\times$ higher than at a $\text{Al}_{0.33}\text{Ga}_{0.67}\text{As}/\text{GaAs}$ interface (≈ 0.81 eV instead of ≈ 0.24 eV). This means that much higher electron densities are needed to make this barrier thin enough for field emission currents. Secondly, Sn normally acts as a donor by occupying Ga sites in the AlGaAs. Therefore it is probably more difficult to dope AlAs with Sn. This is confirmed by the much longer annealing times (≈ 75 s) which are needed to obtain low-resistance ohmic contacts and the values of the contact resistances which are $\sim 10\times$ higher than on conventional material ($T = 4.2$ K). It would be worthwhile to replace the AlAs layers by AlGaAs and to repeat the experiment.

More information about the composition of the contact region should be accessible by SIMS⁶ and TEM⁷. First attempts have however not yet yielded conclusive results. The relief in the contact surface makes SIMS analysis far from trivial. The best way may be to thin the substrate by etching and perform an analysis from the backside of the contact.

The purpose of the thin Ti layer was, besides improving adhesion, to act as a sieve (analogous to processes used for contacts on Si) and to provide in this way a homogeneous Sn diffusion. The resistances of Ti/Sn contacts turn out to be much lower than those of Pd/Sn contacts [28]. Therefore Ti seems to be especially suitable as material. Maybe this is because at room temperature Ti and GaAs react very uniformly, forming a thin amorphous layer with an atomically flat interface [16]. Since 100 nm thick Ti layers are known to form rough interfaces after 20 minutes at 450°C. it is possibly important that our Ti layer is as thin as possible (but continuous).

It can't be excluded that other alloys are formed during annealing, maybe even superconducting with a high critical temperature [29]. However, to make significant remarks about this material analysis is needed first.

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⁶SIMS=secondary ion mass spectroscopy

⁷TEM=transmission electron microscopy

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Chapter 8

Andreev reflection at superconducting contacts to GaAs/AlGaAs heterostructures

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Abstract

Highly transmissive ohmic contacts to the two-dimensional electron gas in GaAs/AlGaAs heterostructures have been made. For these contacts, which are of μm -scale, a newly developed process of Ti/Sn evaporation and diffusion has been used. Devices consisting of these superconducting contacts combined with gate structures have shown clear evidence for the occurrence of Andreev reflection. Moreover a marked effect of the gate voltage on the $dV/dI - V$ -characteristics has been found, which proves that superconductivity has been induced into the semiconductor.

8.1 Introduction

The progress in the microfabrication techniques and the availability of very high-mobility two-dimensional semiconductors have induced recently a renewed interest in the realization and the understanding of the Josephson field-effect transistor or superconducting transistor [1]. This device is based on the proximity coupling [2] of two superconducting electrodes (S), connected by a semiconductor (N). The supercurrent flowing between the electrodes is dependent on the electron density in the semiconductor, which can be varied by a gate voltage.

Until now all realized devices [3–8] are operating in the dirty limit, where the electron transport is diffusive [9]. In the clean limit, where the electron transport is ballistic (at least over the normal coherence length ξ_N), different and new effects are expected [10–13], such as the quantization of the critical supercurrent in a superconducting quantum point contact [14, 15]. GaAs/AlGaAs heterostructures are very suitable to study these kind of effects, because of the very high mobility of the two-dimensional electron gas (2DEG) and the excellent gating properties. In this paper we will discuss the method and the first results [16] of the study of electron transport in such a superconductor/2DEG system.

The superconducting properties of the SNS system critically rely on the transport characteristics of the SN interface. In buried 2DEG systems the length of the contact region approximately equals the distance of the 2DEG to the surface, in our heterostructures $d \approx 80$ nm. In this region phase coherence has to be preserved in order to induce superconductivity into the 2DEG. This requires that the normal coherence length ξ_N has to be at least of the same order of magnitude as d .

If phase coherence is not preserved over the distance L between the two contacts there will be no supercurrent through the device. However, if the contact is highly transmissive, i.e. barrierless, another effect can still arise: Andreev reflection [17, 18]. This process, where an electron from the N region is retroreflected as a hole (and vice versa) with a Cooper pair being injected into the superconductor (or the reverse process), gives rise to an excess current and reduces the resistance [19]. The probability of Andreev reflection depends strongly on the transmission of the NS interface [20]. In particular at interfaces with semiconductors the Schottky barrier can introduce normal reflection and reduce this probability dramatically. This stresses the importance of the contact properties for the operation of the device.

8.2 Fabrication of the samples

Reducing the Schottky barrier can be accomplished by heavily doping the interface layer. Tin is the only element that forms ohmic contacts to GaAs/AlGaAs by doping and which has a reasonable superconducting critical temperature T_c . While it is no problem to make ohmic contacts by thermally diffusing macroscopic Sn pieces into the heterostructure, it is by no means trivial to contact the 2DEG by evaporated Sn film microstructures. As soon as the temperature exceeds the melting point (232 °C), the films tend to contract into droplets on the surface, without significant diffusion into the semiconductor. To suppress this, we inserted a thin layer of a metal that improves the wetting of the surface and that has a high melting point. This layer acts as a sieve during the annealing and helps to form a homogeneous contact. Although several metals (Ni, Pd, Pt, etc.) meet these requirements, we use Ti, because it is also a superconductor itself ($T_c \approx 0.4$ K).

To obtain reliable diffusion the GaAs surface has to be treated before the metal evaporation. Immediately before a sample with the contact pattern in PMMA-resist is mounted in the evaporation chamber, the oxides are removed by etching for 1 minute in 1:25 $\text{NH}_4\text{OH}(25\%):\text{H}_2\text{O}$. When the pressure in the chamber is below 10^{-7} Torr, the Ti source is preheated to clean its surface. Subsequently 10 nm Ti is evaporated with a high rate (≈ 1 nm/s) to avoid oxidation. This is followed by a 200 nm Sn layer (also after preheating of the source). After lift-off in acetone the films are annealed during 5 s at 405 °C in a nitrogen-atmosphere. Increased annealing times did not result in lower contact resistances.

The Ti/Sn contacts realized in this way show a very low resistance of less than $\approx 10 \Omega$ at $T = 80$ mK for a contact of width $W = 10 \mu\text{m}$. Note that this resistance is of the same order as the two-dimensional Sharvin resistance [21], $R_{\text{Sh}} = \frac{h}{2e^2} \cdot \frac{\lambda_F}{2W}$ (λ_F the Fermi wavelength), indicating a small resistance contribution of the transition region. To study the superconducting properties in more detail, devices with a layout as in fig. 1 are fabricated. The centre consists of two rectangular Ti/Sn contacts at a distance of $L = 1 \mu\text{m}$. These contacts are placed on an etched mesa for device isolation. Each contact is connected to bond pads by two Ti/Au leads. Moreover two gates are situated at the sides of the superconducting contacts to confine the region where the current can flow. Four larger Ti/Sn contacts are added to enable tests (e.g. quantum Hall measurement to test two-dimensionality).

Figure 8.1: *Layout of the centre of the sample on scale; the dashed lines are the boundaries of the etched mesa, the black areas are gold layers, the two rectangles are the Ti/Sn contacts*

The GaAs/AlGaAs heterostructure was grown by molecular beam epitaxy. The 2DEG has a mobility $\mu_e = 42 \text{ m}^2/\text{Vs}$ and an electron density $n_e \approx 3.1 \times 10^{15} \text{ m}^{-2}$, which results in an elastic mean free path $l_e = 4 \text{ }\mu\text{m}$.

For all processing steps (except the mesa-etching) we have used electron beam lithography. To align all the etching and metallization steps we needed gold markers, which were fabricated also by electron lithography in a separate processing step. The achieved relative alignment of the patterns, which was done by automatic marker search, is better than $\approx 20 \text{ nm}$.

8.3 Results of the measurements

The measurements were performed in a dilution refrigerator at temperatures down to $\approx 80 \text{ mK}$, employing a current-biased ac lock-in technique as well as dc methods.

First the differential resistance dV/dI between the two small Ti/Sn

contacts of the device (at zero gate voltage V_g) has been measured as a function of the bias voltage V at different temperatures. Because of bonding problems the measurements on this sample were done in a three-terminal configuration. The resulting series resistance (due to a lead) is estimated to be 100-150 Ω . The results are shown in figure 2.

Figure 8.2: The differential resistance as a function of the voltage at different temperatures ($V_g = 0$)

At low temperatures a low resistance region exists around zero-voltage, flanked by two peaks. As we will explain in the following, this is clear evidence for Andreev reflection at a highly transmissive SN interface.

Employing the larger contacts it is also possible to measure the resistance of a single small contact. In this way it was found that a comparable structure occurs in the $dV/dI - V$ -characteristics for each individual small contact, underlining that this is not only an unintentional feature of just one of the small contacts. The measured shape is typical for Andreev reflection in clean NS structures [22]. The fact that the bottom of the low resistance region is very flat and does not show a peak around $V = 0$ V implies that the transmission of the NS interface is very close to 1. Since no indication of multiple Andreev reflection is visible, the two NS interfaces seem to be independent. The voltage where the step in the resistance occurs should correspond with $2\Delta/e$. In the

case of full Andreev reflection the contact resistance will double [19, 20] at $V > 2\Delta/e$. From the measured resistance step at 80 mK we find a series resistance of 140 Ω and a (normal state) contact resistance of 9 Ω . The voltage at the resistance step yields a superconducting gap of ≈ 10 μeV . The typical shape of the curve disappears completely between 0.2 K and 0.5 K. This is in agreement with the critical temperature of bulk Ti ($T_c \approx 0.4$ K), suggesting that Ti could be the relevant superconductor at which interface the Andreev reflection takes place. The value found for the superconducting gap is nevertheless considerably smaller than the BCS-value $2\Delta = 3.5k_B T_c$. Therefore, one could think about a quantum interference effect [20, 23], for which the relation between the voltage and the temperature at which the effect is suppressed is given by $eV_0 = k_B T_0$. However, it is very difficult to imagine how quantum interference could produce such a steep step in the $dV/dI - V$ -curve. This is also the problem in comparing our results with the experiment of ref. 24.

Second, the influence of the gate voltage on the $dV/dI - V$ -characteristics has been studied. As can be seen in figure 3 the resistance dip disappears completely at large negative gate voltages ($V_g \lesssim -1$ V). We have checked that this effect is not caused by leakage currents from the gates. While the height of the dip reduces rather abruptly to zero, the width of the flat, low resistance region decreases gradually. This underlines that this feature is associated with the superconducting gap. Although this effect is not understood yet, it proves that superconductivity has been induced into the semiconductor.

8.4 Conclusions

In summary, we have developed the first method for making low-resistance superconducting contacts to the 2DEG in GaAs/AlGaAs heterostructures that are sufficiently transmissive to observe Andreev reflection. This process uses a thin intermediate metal layer (Ti) between Sn and the heterostructure, which improves the wetting of the surface and can act as a sieve during the annealing. In this way superconducting contacts of μm -scale have been made in combination with gate structures. The temperature and gate-voltage dependences of the $dV/dI - V$ -curves show that Andreev reflection takes place in the sample and that superconductivity has been induced into the semiconductor. The flatness of the low resistance region in the $dV/dI - V$ -curve shows that the transmis-

Figure 8.3: *The differential resistance as a function of the voltage at different gate voltages ($T \approx 80$ mK), resp. 0 mV, -300 mV, -400 mV, -600 mV, -1000 mV, -1500 mV and -2000 mV*

sion of the contact is very close to 1. Therefore these superconducting contacts seem to be very suitable to study ballistic SNS systems.

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Chapter 9

Mesoscopic phenomena in superconductor/2DEG systems

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9.1 Introduction

In the last decade a new branch of physics arose, studying systems with sizes between microscopic and macroscopic dimensions: mesoscopic physics. This field has yielded many interesting results [1], as these proceedings also show. The transport properties of low-dimensional systems are, in fact per definition, strongly determined by the boundaries. Whether we consider a quantum point contact, a quantum wire or a quantum dot, always the confinement is essential. Therefore it is obvious that new effects will be encountered, if some of the boundaries

can be formed by superconductors. Due to Andreev reflection [2] the scattering of electrons at the boundaries will change dramatically.

The aim of our research is to study the interplay of Andreev reflection and the mesoscopic effects which are known to occur in semiconductor nanostructures. We want to use semiconductors, so that the normal N region can be patterned by means of gate structures, and since we want to study phase-coherent effects in particular, we need a semiconductor with a high electron mobility. That is why the use of a two-dimensional electron gas (2DEG) is preferred.

In the past several superconductor/2DEG systems have already been studied: a Si MOSFET [3,4], a native inversion layer on *p*-type InAs [5], an accumulation layer in $\text{Al}_{0.5}\text{Ga}_{0.5}\text{As}/\text{GaAs}$ [6], an InAs-AlSb quantum well [7] and $\text{GaAs}/\text{Al}_{0.33}\text{Ga}_{0.67}\text{As}$ heterostructures [8,9]. Each of these material systems has its own advantages and disadvantages. While it is relatively easy to make contact between InAs and a superconductor because of the absence of a Schottky barrier, it is certainly not trivial to make good gate structures on this material. In Si MOSFET's supercurrents have indeed been observed and could be modulated by a gate; however, the phase coherence length is so small, that the distance between the superconducting contacts is not enough to place split gate patterns.

We have chosen to use $\text{GaAs}/\text{Al}_{0.33}\text{Ga}_{0.67}\text{As}$ heterostructures, because of its high mobility ($\mu \sim 10^2 \text{ m}^2/\text{Vs}$) and its excellent gating properties. The large elastic mean free path ($l_e \sim 10^{-2} \text{ mm}$) gives the possibility to enter the ballistic transport regime, while the distance between the superconducting contacts can be large enough for gate structures like a quantum point contact or a quantum dot. One of the effects which are theoretically predicted to occur in such systems, is the quantization of the critical supercurrent [10,11]. However, until now supercurrents in a ballistic superconductor/2DEG/superconductor structure that can be controlled by a gate have not yet been realized. On the other hand, new effects have been observed due to phase-coherent Andreev reflection in mesoscopic systems. This will be the subject of the following, by discussing our experiments on $\text{GaAs}/\text{AlGaAs}$ heterostructures (section 9.4). Before that we will shortly describe the most important theoretical concepts (section 9.2) and the fabrication process of our samples (section 9.3).

Figure 9.1: Schematic view of an Andreev reflection process

9.2 Theoretical concepts

If an electron from a normal material with an excitation energy E below the superconducting gap Δ reaches the interface with a superconductor (S), it can't enter the superconductor because there are no states available (see figure 9.1). Andreev [2] found out that the electron can be reflected as a hole, while a Cooper pair goes on in the S region. While in normal reflection the angle of incidence equals the reflection angle, in an Andreev reflection the hole traces back the path of the incoming electron (the time-reversed path) [12]. It is evident that this retroreflection feature has a large influence on quantum interference. The second important feature of Andreev reflection is the already mentioned particle conjugation. As can be seen in figure 9.1 we end up with a Cooper pair (charge $2e$) in the superconductor, while we started with an electron (e) in the N region. This means that Andreev reflection effectively doubles the current; experimentally this will be observed as an increase of the conductance.

However, if the NS interface is not ideal, the probability for Andreev reflection will dramatically decrease in the case of a small barrier. Since Andreev reflection is a two-particle process in contrast to normal electron tunneling, the conductance then will decrease at the transition from an NN interface to an NS interface.

So on beforehand we cannot say whether the conductance of the interface will go up or down at the superconducting transition. The nett

effect depends strongly on the transmission coefficient of the interface [13].

For the study of mesoscopic effects it is useful to express the conductance of the NS interface G_{NS} in terms of the eigenvalues T_n of the (normal-state) transmission matrix product $tt^\dagger$, like Landauer [14] did for the normal state:

$$G_{NN} = \frac{2e^2}{h} \sum_{n=1}^N T_n. \quad (9.1)$$

Beenakker [15] recently found a generalization of this formula for the NS interface ($B = 0$ T):

$$G_{NS} = \frac{4e^2}{h} \sum_{n=1}^N \frac{T_n^2}{(2 - T_n)^2}. \quad (9.2)$$

If this expression is applied to a ballistic quantum point contact (that means either $T_n = 0$ or $T_n = 1$) near an NS interface, it is clear that the conductance will be quantized in units of $4e^2/h$ instead of the normal value of $2e^2/h$.

Now we consider a more specific case, namely an NS interface with a disordered N region. Here “disordered” means that the elastic mean free path l_e is much smaller than the length L of the N region. The transmission of an eventual barrier at the NS interface is called Γ . This system has recently theoretically been described by Marmorkos *et al.* [16] and is schematically shown in figure 9.2. Classically (i. e. matching of currents instead of amplitudes) the resistance in respectively the superconducting and normal state is given by:

$$R_{NS}^{\text{class}} = \frac{h}{2e^2N} \left(T_d^{-1} + 2(1 - \Gamma)\Gamma^{-2} \right) \quad (9.3)$$

and

$$R_N^{\text{class}} = \frac{h}{2e^2N} \left(T_d^{-1} + (1 - \Gamma)\Gamma^{-1} \right) \quad (9.4)$$

with T_d the transmission of the disordered region. Looking at these formulas we can distinguish two limits:

- high barrier ($\Gamma \ll 1$): $R_{NS}^{\text{class}} \gg R_N^{\text{class}}$, because the contribution of the barrier is $\propto \Gamma^{-2}$ instead of $\propto \Gamma^{-1}$
- without barrier ($\Gamma = 1$): $R_{NS}^{\text{class}} = R_N^{\text{class}}$.

Figure 9.2: *Schematic overview of an NS interface with disordered N region*

However for both limits deviations occur, when quantum interference is taken into account [15–17]:

- high barrier ($l_e/L \ll \Gamma \ll 1$): $R_{NS} \approx R_N^{\text{class}} \ll R_{NS}^{\text{class}}$. This is called reflectionless tunneling, because it is as if the Andreev-reflected hole is not reflected by the barrier. Experimentally this will result in a conductance peak
- without barrier ($\Gamma = 1$): $R_{NS} > R_{NS}^{\text{class}}$ due to an enhancement of weak localization: it is as if Andreev reflection effectively doubles the length of the disordered region, which will show up in measurements as an conductance dip.

Experimental observations of three signatures of (phase-coherent) Andreev reflection will be discussed in this paper: conductance quantization in enhanced units, enhancement of weak-localization effects and reflectionless tunneling.

9.3 Sample fabrication

For most samples a “conventional” GaAs/AlGaAs heterostructure has been used. This material is grown by Molecular Beam Epitaxy and consists of a GaAs substrate with a $4\text{ }\mu\text{m}$ undoped GaAs layer, a 40 nm undoped $\text{Al}_{0.33}\text{Ga}_{0.67}\text{As}$ spacer, a 38 nm n -doped $\text{Al}_{0.33}\text{Ga}_{0.67}\text{As}$

Figure 9.3: *Side-view of the centre of the sample*

layer and a 17 nm undoped GaAs caplayer. The electron mobility is $\mu \approx 110 \text{ m}^2/\text{Vs}$ and the electron density is $n_e \approx 2.4 \times 10^{15} \text{ m}^{-2}$ (at $T \approx 5 \text{ K}$).

The major problem is the coupling between a superconductor and the buried 2DEG. We have developed a diffusion process for making highly transmissive superconducting Sn/Ti-contacts, which have shown clear evidence of Andreev reflection [8,18]. Using this process we have made a sample with a layout as shown in figure 9.3 on this heterostructure. The centre of the sample consists of two $10 \mu\text{m}$ wide superconducting contacts at a mutual distance of 800 nm. Each of these contacts is attached to two large bond pads for the connection to the outside world. Splitgates forming a quantum point contact (width $\approx 150 \text{ nm}$, length $\approx 90 \text{ nm}$) are positioned between the small contacts. Outside the contacts the gates are widened, so that they can confine the current flow to the region between the small superconducting contacts. The whole sample is surrounded by a wide gate (a moat gate) for device isolation.

The patterning of all three layers (contacts, wide gate and quantum point contact) is done by 100 keV electron beam lithography in a double PMMA resist layer to reduce negative influences of the lithographic proximity effect. The alignment of these three lithography steps is achieved

by automatic marker search on especially prepared gold markers (with an inaccuracy below 30 nm). With this process also six thin gates forming a quantum dot have been fabricated between superconducting contacts (distance ≈ 900 nm) on a test wafer, proving the reliability and flexibility of this nanolithography and alignment technique.

The measurements of reflectionless tunneling have been performed on a different heterostructure with a “shallow” 2DEG. For this material a different layer structure has been grown on a GaAs substrate: respectively, a $0.5\ \mu\text{m}$ undoped GaAs layer, a 10 nm undoped AlAs layer, a δ -doped 0.56 nm GaAs layer sandwiched between two 1.13 nm $\text{Al}_{0.33}\text{Ga}_{0.67}\text{As}$ layers, again a 10 nm undoped AlAs layer and a 5.4 nm undoped GaAs caplayer. In this way the 2DEG (with $\mu \approx 60\ \text{m}^2/\text{Vs}$ and $n_e \approx 3.6 \times 10^{15}\ \text{m}^{-2}$) is situated only ≈ 28 nm below the surface. On this material $200\ \mu\text{m} \times 200\ \mu\text{m}$ superconducting contacts are placed on a $200\ \mu\text{m}$ wide etched mesa at several distances (a so-called TLM structure).

The resistance measurements on the samples were performed in a dilution refrigerator, employing a current-biased ac lock-in technique as well as dc methods.

9.4 Results of the measurements

9.4.1 Conductance quantization in enhanced units

In figure 9.4 the conductance of the sample with the quantum point contact is presented as a function of the gate voltage V_g ($T = 10\ \text{mK}$). It can be clearly seen that at zero magnetic field the conductance is quantized in units larger than $2e^2/h$. We think this can only be due to Andreev reflection. In this respect it is important to note that these are the original measurement data, i. e. without a correction for a series resistance.

The fact that the steps are practically equidistant for $B = 0\ \text{T}$ means that there is no significant series resistance. The figure also shows that the effect is suppressed by small magnetic fields; even the remanent field of the superconducting magnet is enough to destroy the enhancement.

By means of formula (9.2) we can make an estimation for an effective, averaged transmission probability T (assuming mode independence). From the fact that the observed quantization unit is $\approx 2.2e^2/h$ instead of the ideal value $4e^2/h$, we find $T \approx 0.85$. To check this we

Figure 9.4: *The differential conductance as a function of the voltage on the quantum point contact gates*

can do an analogous estimation in the normal state using formula (9.1), which indeed yields the same value. This suggests that the transmission of an eventual barrier at the NS interface Γ is larger than 0.85. Following the theory [16] this is in the regime where enhanced weak-localization effects can be expected. Indeed this has been observed as will be discussed in the next section.

9.4.2 Enhanced weak-localization effects

At zero gate voltage the conductance shows a sharp minimum at zero bias, which is rapidly suppressed by a magnetic field (figure 9.5) and increased temperature (figure 9.6). We think this minimum consists of two superimposed phenomena: a broad minimum due to the opening of the superconducting gap at a non-ideal NS interface [13] and a sharp dip caused by enhanced weak localization [15]. This hypothesis is confirmed by measurements on a sample in which the major part of the thin gate between the superconducting contacts doesn't function, but in which the current flow still can be confined by means of the wider gates at the sides of the contacts (figures 9.7 and 9.8). Figure 9.7 clearly shows

Figure 9.5: The normalized differential conductance as a function of the bias voltage for several magnetic fields ($T = 10$ mK)

Figure 9.6: The differential conductance as a function of the bias voltage at several temperatures ($B = 0$ mT)

Figure 9.7: The differential conductance at $V_g = -400$ mV as a function of the bias voltage at several temperatures ($B = 0$ mT)

Figure 9.8: The differential conductance at $V_g = -400$ mV as a function of the bias voltage for several magnetic fields ($T = 10$ mK)

that at temperatures below 0.9 K a small dip arises superimposed on the broad minimum. This dip also vanishes in a small magnetic field (figure 9.8) which is in agreement with the fact that weak localization is suppressed when time reversal symmetry is broken. The depth of the dip is also in good agreement with theory which predicts a magnitude in the order of e^2/h [19]. The disorder which causes the weak localization is probably introduced near the contacts by the diffusion process.

Also in the original sample (figures 9.5 and 9.6) the sharp dip is suppressed at $B \approx 35$ mT and $T \approx 0.75$ K. It is difficult to predict the

voltage at which the enhanced weak localization is suppressed, since the Fermi velocity and the elastic mean free path of the disordered region are not known. However, the threshold voltage and the threshold temperature should be related by $eV_{\text{th}} \approx 4k_{\text{B}}T_{\text{th}}$. The observed values are in good agreement with this prediction ($V_{\text{th}} \approx 0.25$ mV).

The enhanced weak localization can however also be suppressed by a large negative voltage on the nearby gates. At $V_{\text{g}} = -890$ mV not a conductance minimum but a maximum is found around zero bias and also the sign of the temperature dependence is opposite [20, 21]. We think the suppression is caused by the formation of a potential barrier in the quantum point contact; this lowers the transmission of the sample, which is then not in the regime of enhanced weak localization anymore. Anyway, it proves that phase coherence is preserved into the 2DEG.

From the theoretical expression for the threshold magnetic field [16] $B_{\text{th}} = h/(el_{\varphi}^2)$ we find an estimation for the phase coherence length l_{φ} of the order of $0.4 \mu\text{m}$. This is in agreement with the observed gate influence.

9.4.3 Reflectionless tunneling

Reflectionless tunneling was the first signature of phase-coherent Andreev reflection that was experimentally observed. It was measured for the first time in a Nb/InGaAs junction by Kastalsky *et al.* [22]. Later it has been observed in other material systems [4, 7, 21, 23]. Van Wees [24] was the first to realize that the measured conductance peak could be explained by phase-coherent Andreev reflection; this concept was elaborated by others [15, 16].

We have observed reflectionless tunneling most clearly in samples on the shallow 2DEG. The data are shown in the figures 9.9 and 9.10.

Below the critical temperature of tin ($T_{\text{c}} = 3.7$ K for bulk) a broad conductance minimum arises because of “classical” Andreev reflection at an NS interface with a non-unity transmission [13]. At very low temperatures ($T < 0.3$ K), however, a small conductance peak appears, due to phase-coherent Andreev reflection (figure 9.9). This temperature is in correspondence with the voltage at which the peak vanishes. The peak is also suppressed by a small magnetic field of $B \approx 20$ mT (figure 9.10). This suggests that $l_{\varphi} \approx 0.45 \mu\text{m}$, which means that phase coherence is preserved into the 2DEG.

The fact that we observe reflectionless tunneling in this shallow

Figure 9.9: *The differential conductance as a function of the bias current at several temperatures*

2DEG means that the NS interface is less transmissive as in the conventional 2DEG. This is in agreement with the higher contact resistances we find on this material, probably due to the different layer structure.

9.5 Conclusions

We have reported the experimental observation of three signatures of phase-coherent Andreev reflection in superconductor/2DEG systems. The conductance of a quantum point contact near a superconductor is quantized in units larger than $2e^2/h$. At zero gate voltage an enhancement of weak localization by Andreev reflection has been observed for the first time. In a different heterostructure reflectionless tunneling has been measured for the first time in GaAs/AlGaAs.

The fact that in these effects phase coherence is preserved into the 2DEG shows that GaAs/AlGaAs heterostructures are very promising for the study of the interplay between Andreev reflection and mesoscopics.

Figure 9.10: *The differential conductance as a function of the bias voltage for several magnetic fields*

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Chapter 10

New experimental signatures of phase-coherent Andreev reflection

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Abstract

We report the first experimental observation of a new signature of phase-coherent Andreev reflection. In highly transmissive, superconducting contacts to the high mobility two-dimensional electron gas in GaAs/AlGaAs heterostructures, a sharp conductance dip is found around zero bias voltage, which is ascribed to an enhancement of weak localization by Andreev reflection. This new phenomenon is not only suppressed by increased temperature, a small magnetic field or a large bias current, but also by applying a large negative voltage to a nearby gate. This behaviour is in agreement with recent theoretical predictions.

10.1 Introduction

Interference effects in normal metals or semiconductors (N) will be influenced dramatically if a superconductor (S) is brought into the vicinity. Due to Andreev reflection [1] an electron can be retroreflected as a hole at the NS interface (and v.v.). In particular within a distance of the phase-coherence length l_φ from the superconductor this will have large consequences for the electronic transport properties.

One of the earliest experimental observations of phase-coherent Andreev reflection was done a few years ago [2]. In the $dI/dV - V$ -curve of a Nb/InGaAs junction a sharp conductance peak was found around zero bias. This feature was described in terms of phase-coherent Andreev reflection [3–5]. Later this phenomenon, now known as reflectionless tunneling, has also been observed in other materials [6–9]. It was realized [4,5] that other new interference effects can be expected if mesoscopic systems are combined with superconductors. In this Letter we will discuss the first experimental observation of such a new phenomenon: the enhancement of weak localization.

Semiconductors have the advantage above metals that the electron density can be modulated by means of gate structures. Especially in mesoscopic physics nanostructures are often fabricated by means of fine gate patterns [10]. It is obviously a big advantage if this valuable experience in submicron technology can be used for the study of phase-coherent Andreev reflection. Therefore we developed a process for making diffused superconducting contacts to the two-dimensional electron gas (2DEG) in GaAs/AlGaAs heterostructures [11]. Clear indications of Andreev reflection have been found in these contacts and the transmission of the NS interface turned out to be close to unity [12].

We used this process to make a sample with a quantum point contact between two 10 μm wide superconducting contacts at a mutual distance of $\approx 0.8 \mu\text{m}$. The 2DEG has an electron density $n_e \approx 2.4 \times 10^{15} \text{ m}^{-2}$ and a mobility $\mu \approx 110 \text{ m}^2/\text{Vs}$. This yields an elastic mean free path $l_e \approx 9 \mu\text{m}$. The measurements are performed in a dilution refrigerator with a base temperature of $\approx 10 \text{ mK}$, using a current biased lock-in technique.

Measurements of the conductance of this sample as a function of the voltage V_g on the gates which form the quantum point contact have shown conductance quantization in steps larger than $2e^2/h$ [13], providing evidence for Andreev reflection. From the observation that the

quantization unit is approximately $2.2e^2/h$ instead of $4e^2/h$ (which is expected for the ideal, clean case) we can make a rough estimation of the effective transmission probability of the sample. Under the assumption that the transmission T is mode-independent we find an effective value of 0.85, using the generalization of the Landauer-Büttiker formula for an NS junction [4]:

$$G_{\text{NS}} = \frac{4e^2}{h} \sum_{n=1}^N \frac{T_n^2}{(2 - T_n)^2}. \quad (10.1)$$

This effective transmission value is also consistent with an analogous estimation using the Landauer-Büttiker formula in the normal state. We think, that the transmission is reduced due to disorder (induced by the diffusion process) near the NS interface and a non-unity transmission of the interface itself. Anyway, this argument shows that the transmission of an eventual barrier at the NS interface is larger than ≈ 0.85 , which is clearly in the regime where enhanced weak localization effects can be expected, according to [5].

10.2 Results and discussion of the measurements

In figure 10.1 the differential conductance of the sample is shown as a function of the bias voltage at zero gate voltage V_g . Without magnetic field a broad minimum with a sharp dip is observed around zero bias voltage. A similar $dI/dV - V$ -curve is also found for contacts which are $470 \mu\text{m}$ apart, so this phenomenon must be related to separate NS interfaces and not to an coupled SNS system. If the magnetic field is increased, the whole feature vanishes at a value well below 60 mT.

The effect of increasing temperature can be seen in figure 10.2: after a rapid rounding of the minimum the whole structure has disappeared at 4.2 K. Detailed inspection of the low-temperature curve (without magnetic field) suggests that the minimum consists of two superimposed phenomenon. This suggestion is confirmed by measurements on a sample where the major part of the gates between the two small superconducting contacts doesn't function, but the current path still can be confined by the wider gates besides the small contacts. The conductance of this sample is presented in figure 10.3. At $T = 0.9 \text{ K}$ only a broad minimum is observed, while at lower temperatures a sharp conductance dip is

Figure 10.1: *The normalized differential conductance as a function of the bias voltage at several magnetic fields ($V_g = 0$ mV)*

superimposed on this. This clearly shows the two constituents of the effect. In figure 10.2 we see that at $T = 0.9$ K the conductance minimum has become much more rounded than at lower temperatures, while the rest of the curve has barely changed, which is in accordance with the concept of two separate effects.

The broad conductance minimum can be well described by Andreev reflection at a non-ideal NS interface [14]. If the interface has a finite transmission a reduction in the conductance is expected below the superconducting gap, due to the fact that Andreev reflection is a two-particle process, in contrast with electron tunneling which is a one-particle process. The value of the temperature and magnetic field at which the minimum is suppressed are in reasonable agreement with the superconducting properties of tin ($T_c = 3.7$ K and $H_c = 30$ mT for bulk Sn). That the bias voltage at which the conductance reduction in our measurements starts isn't exactly equal to twice the superconducting gap Δ of tin can be explained by a series resistance of $\approx 10 \Omega$. Of course it is also possible that the superconducting properties have been changed by impurities or even alloys in the tin.

We attribute the small, sharp conductance dip to phase-coherent

Figure 10.2: The differential conductance as a function of the bias voltage at several temperatures ($V_g = 0\text{ mV}$)

Andreev reflection. In a junction between a superconductor and a disordered normal material weak localization effects can be enhanced by Andreev reflection, if the barrier at the NS interface is not too high [4,5]. It is as if Andreev reflection effectively doubles the length of the disordered region. In contrast to conventional weak localization effects this phenomenon changes the (low) V -dependence of the conductance: around zero bias the conductance will be reduced, but at larger bias voltages phase coherence between the electrons and Andreev-reflected holes is broken and the conductance approaches again its normal value. This is what we see in figures 10.2 and 10.3; the sharp dip is suppressed at a critical voltage $V_c \approx 0.25\text{ mV}$. Evidently, weak localization effects can be suppressed by a small magnetic field, since then time reversal symmetry is broken. This is also observed in the measurements (figure 10.1 and 10.4). The conductance first increases as a function of the magnetic field, because weak localization is suppressed, before it decreases due to the square magnetoresistance of the 2DEG. In $dI/dV - V$ -curves in magnetic fields it is noticed that at the same B -value the small conductance dip has practically vanished.

Figure 10.3: The differential conductance as a function of the bias voltage at several temperatures ($V_g = -400$ mV)

Theoretically [5] the critical magnetic field is given by

$$B_c = \frac{h}{el_\varphi^2}, \quad (10.2)$$

if the phase-coherence length is smaller than the length of the disordered normal region. From this equation we estimate l_φ to be of the order of $0.4 \mu\text{m}$. A quantitative prediction of V_c by means of the expression

$$eV_c = \frac{\pi \hbar v_F l_e}{2 l_\varphi^2} \quad (10.3)$$

is difficult since the Fermi velocity v_F and the mean free path l_e of the disordered region are not known. However the experimentally found V_c and T_c should be related by $eV_c = 4k_B T_c$. From $V_c \approx 0.25$ mV we would expect $T_c \approx 0.75$ K. This is in agreement with the measurements. Moreover the absolute depth of the conductance dip should be of the order of e^2/h for a single interface. Indeed the experimental observation for $V_g = -400$ mV, where the current path is confined within the 2DEG region between the two superconducting contacts, has the right order of magnitude [15].

Figure 10.4: The differential conductance as a function of the magnetic field ($V_g = -400$ mV)

The temperature dependence of the normalized differential conductance change is plotted in figure 10.5. At low temperatures and without magnetic field applied the conductance is increasing with temperature. If the enhanced weak localization is suppressed, a very different dependence is observed; even the sign of the conductance change is opposite. As can be seen in the graph this suppression can be achieved by applying a small magnetic field or a large DC current [16]. The most important observation, however, is the fact that the enhanced weak localization can also be suppressed by a large negative voltage on the nearby gates. This proves that phase coherence is preserved down to in the 2DEG. This is also predicted by the estimation of $l_\varphi \approx 0.4$ μm , which is much larger than the distance of the 2DEG from the surface (≈ 85 nm) and, evidently, also follows from the observed conductance quantization in enhanced units. We think that the explanation for the suppression is the formation of a potential barrier in the quantum point contact. This lowers the total transmission of the sample, which is then not in the regime of enhanced weak localization anymore, but rather in the regime of reflectionless tunneling [5]. Indeed we have observed a maximum in the $dI/dV - V$ -curve at zero bias for $V_c = -890$ mV, instead of the

Figure 10.5: *The temperature dependence of the normalized conductance change*

minimum found for $V_g = 0$ mV.

10.3 Summary

In summary, we have presented the first experimental observation of the enhancement of weak localization by Andreev reflection. In correspondence with theoretical predictions the effect is suppressed by increasing temperature, a small magnetic field or a large DC current. Moreover the enhancement vanishes, when a large negative voltage is applied to the nearby quantum point contact gates. This proves that phase coherence between Andreev-reflected particles exists in the 2DEG.

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- [16] Since the experiment is current-biased a large DC current is applied. Physically, however, Andreev reflection (and thus the enhanced weak localization) is suppressed, because the resulting voltage over the NS interface is much higher than the superconducting gap Δ .

Summary

In many fields of physics interest is growing for low-dimensional systems. On the one side because they can sometimes be described by simpler models than three-dimensional systems, on the other hand because new fundamental phenomena are encountered here. In this thesis electron transport has been studied in systems that combine low-dimensional semiconductor structures with superconducting contacts.

The starting-point is the two-dimensional electron gas (2DEG) in GaAs/AlGaAs heterostructures. By means of nanolithographically patterned metal electrodes (gates) on top of this material, the electrons can be confined also within the plane of the 2DEG. The thus obtained mesoscopic systems have already shown a wealth of interesting phenomena, like conductance quantization in a quantum point contact and discretization of energy levels in a quantum dot.

In the first part of this thesis an example of an intriguing semiconductor structure, fabricated by state-of-the-art electron beam lithography, has been discussed: by means of a grid-shaped gate a two-dimensional periodic potential (a 2D superlattice) is induced in the underlying 2DEG. The strength of this potential can be modulated by varying the gate voltage. In this way different transport regimes can be obtained: from a weak periodic potential via an artificial crystal to an array of weakly coupled quantum dots. Because of the very small dimensions of the grid phase coherence is preserved, which gives the possibility to measure, thus far barely studied, effects of two-dimensional coupling.

In low magnetic fields (chapter 3) minima in the magnetoconductance of the sample have been observed, due to commensurability of cyclotron orbits with the lattice periodicity. These minima correspond with localized orbits. Superimposed on the minima small oscillations periodic in the field have been observed. For the first time they have been found in several minima. A study of their periods indicates that the localized orbits are phase-coherently coupled all over the grid.

In high magnetic fields (chapter 4) periodic oscillations in the magnetoconductance have been found on the positions of the Hall plateaux. Comparing grids with different numbers of cells shows that the period scales with this number, which suggests two-dimensional coupling of quantum dots, corresponding with band-structure effects in an artificial 2D crystal. Indeed also the magnitude of the period is in good agreement with this description.

In chapter 5 the subject was changed to the combination of the 2DEG in GaAs/AlGaAs with superconductors. After a discussion of theoretical models of interfaces between superconductors and normal materials, a generalized transmission formalism was described and applied numerically in chapter 6.

The realization of superconducting contacts to the 2DEG was the subject of chapter 7. This is far from trivial, since the 2DEG is situated typically 100 nm below the surface of the heterostructure. Therefore a new process for the fabrication of *diffused* superconducting contacts has been developed. Five different metallizations have been characterized by electrical measurements at room temperature and low temperatures, and compared with the normal Ni/AuGe/Ni/Au contact as a reference. The Sn/Ti metallization turned out to provide the most suitable superconducting contact.

The application of this new type of contact in a device with gates was presented in chapter 8. Clear evidence for Andreev reflection has been observed, showing that the interface between the superconductor and the normal region has a very high transmission probability. The remarkable influence of the gate proves that superconducting effects do not only take place in the top layers of the contact, but that the 2DEG plays an essential role.

The last chapters dealt with new mesoscopic phenomena which have been observed in a sophisticated sample, consisting of superconducting contacts near a quantum point contact. Mesoscopic effects were drastically influenced by Andreev reflection. For the first time conductance quantization in enhanced units (i.e. $> 2e^2/h$) has been observed (chapter 9). This again proves that phase coherence is preserved from the superconductor down to the 2DEG. In the same chapter also the first observation of reflectionless tunneling in GaAs/AlGaAs was presented.

Finally, measurements of a new phenomenon caused by phase-coherent Andreev reflection were described in chapter 10: the enhancement of weak localization.

Samenvatting

In steeds meer gebieden van de natuurkunde groeit de belangstelling voor laagdimensionale systemen. Enerzijds omdat deze soms beschreven kunnen worden met simpeler modellen, anderzijds vanwege de nieuwe fundamentele verschijnselen die hierin optreden. In dit proefschrift is elektronentransport bestudeerd in systemen die laagdimensionale halfgeleiderstructuren combineren met supergeleidende contacten.

Het uitgangspunt vormt het tweedimensionale elektronengas (2DEG) in GaAs/AlGaAs-heterostructuren. Door middel van nanolithografisch gestructureerde metalen elektroden (gates) op het oppervlak van dit materiaal kan de beweging van de elektronen ook binnen het vlak van het 2DEG beperkt worden. De zo verkregen mesoscopische systemen hebben al een grote rijkdom aan interessante verschijnselen getoond, zoals geleidingskwantisatie in een quantum puntcontact en discretisatie van energieniveaus in een quantum dot.

In het eerste deel van dit proefschrift is een voorbeeld van een intrigerende halfgeleiderstructuur, gemaakt met state-of-the-art elektronenbundellithografie, besproken: door middel van een roostervormige gate wordt een tweedimensionale periodieke potentiaal geïnduceerd in het onderliggende 2DEG. De sterkte van deze potentiaal kan gemoduleerd worden door de gatespanning te variëren. Op deze manier kunnen verschillende transportregimes verkregen worden: van een zwakke periodieke potentiaal via een kunstmatig kristal naar een netwerk van zwak gekoppelde dots. Vanwege de zeer kleine afmetingen van het rooster is fasecoherentie behouden, hetgeen de mogelijkheid biedt, tot nu toe nauwelijks bestudeerde, effecten van tweedimensionale koppeling te meten.

In zwakke magneetvelden (hoofdstuk 3) zijn minima in de magnetogeleiding waargenomen, veroorzaakt door commensurabiliteit van cyclotronbanen met de roosterperiodiciteit. Deze minima zijn gerelateerd aan gelokaliseerde elektronenbanen. Gesuperponeerd op de minima zijn

kleine oscillaties waargenomen, die periodiek zijn in het veld. Voor de eerste maal zijn zij nu gevonden in meerdere minima. Een analyse van hun perioden geeft aan dat de gelokaliseerde banen fasecoherent gekoppeld zijn over het hele rooster.

In sterke magneetvelden (hoofdstuk 4) zijn periodieke oscillaties in de magnetogeleiding gevonden op de plaatsen van de Hall-plateaus. Een vergelijking van roosters met verschillende aantallen cellen toont dat de periode schaaft met dit aantal, hetgeen tweedimensionale koppeling van quantum dots suggereert, overeenkomend met bandenstructureffecten in een kunstmatig 2D kristal. Inderdaad is ook de grootte van de periode in goede overeenstemming met deze beschrijving.

In hoofdstuk 5 is overgegaan naar de combinatie van het 2DEG in GaAs/AlGaAs met supergeleiders. Na een discussie van theoretische modellen voor grensvlakken tussen supergeleiders en normale materialen is een gegeneraliseerd transmissieformalisme beschreven en numeriek toegepast in hoofdstuk 6.

De realisatie van supergeleidende contacten naar het 2DEG was het onderwerp van hoofdstuk 7. Dit is verre van triviaal, aangezien het 2DEG typisch 100 nm onder het oppervlak van de heterostructuur gelegen is. Daarom is een nieuw proces voor de fabricage van *gediffundeerde* supergeleidende contacten ontwikkeld. Vijf verschillende metallisaties zijn gekarakteriseerd door elektrische metingen bij kamertemperatuur en lage temperaturen, en vergeleken met het normale Ni/AuGe/Ni/Au-contact als referentie. De Sn/Ti- metallisatie bleek het meest geschikte supergeleidende contact op te leveren.

De toepassing van dit nieuwe type contacten in een device met gates is gepresenteerd in hoofdstuk 8. Duidelijke tekenen van Andreev-reflectie zijn waargenomen, die aantonen dat het grensvlak tussen de supergeleider en het normale gebied een zeer hoge transmissiewaarschijnlijkheid heeft. De opmerkelijke invloed van de gate bewijst dat supergeleidende effecten niet alleen in de bovenste lagen van het contact plaatsvinden, maar dat het 2DEG hierin een essentiële rol speelt.

De laatste hoofdstukken handelden over nieuwe mesoscopische verschijnselen die waargenomen zijn in een geraffineerd preparaat, bestaande uit supergeleidende contacten in de buurt van een quantum puntcontact. Quantum effecten werden drastisch beïnvloed door Andreev-reflectie. Voor de eerste maal is geleidingskwantisatie in vergrootte eenheden (d.w.z. $> 2e^2/h$) waargenomen (hoofdstuk 9). Dit bewijst nog eens dat de fasecoherentie behouden is vanaf de supergeleider tot in het 2DEG.

In hetzelfde hoofdstuk is ook de eerste waarneming van reflectieloos tunnelen in GaAs/AlGaAs gepresenteerd.

Tenslotte zijn metingen van een nieuw verschijnsel, veroorzaakt door fasecoherente Andreev-reflectie, beschreven in hoofdstuk 10: de versterking van zwakke lokalisatie.

Curriculum Vitae

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