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ABSTRACTS OF INVITED LECTURES AND
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$D(E_F)$ for an assumed mobility of $\mu = 50m^2/Vs$. Obviously $D(E_F)$ at finite temperatures ($T = 4K$, solid line and $T = 15K$, dotted line) is not sensitive to the pronounced changes in the miniband structure with n_ϕ. Its effect on the longitudinal conductivity $\sigma_{xx}(E_F)$, however, is dramatic (see right panel, note the different scales). The essential effect derives from the velocity matrixelements: taken between states of the same miniband they are large if the miniband has pronounced dispersion, therefore, the band conductivity is dominant for n_ϕ. For flat minibands ($n_\phi = 9/8$) band conductivity is strongly reduced and $\sigma_{xx}(E_F)$ is essentially due to scattering contributions between different minibands, which gives values an order of magnitude smaller than for $n_\phi = 1$. Changing the temperature from $4K$ (solid lines) to 15K (dotted lines) smoothes the dependence on E_F but does not change the absolute values.

Converting these results to resistivity data we get peaks at integer or half integer n_ϕ, i.e. periodic in B, which persist even at higher temperatures of about $10K$ and for varying electron densities. In this respect they differ from other coexisting structures connected with classical (commensurability peaks) or semiclassical (SdH oscillations) mechanisms, which disappear at higher temperature. We expect that these B periodic quantum peaks whose origin is in the magnetic miniband structure should be observable in properly designed lateral superlattices [12].

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References

1. J. Zak, Solid State Physics Vol. 27, eds. H. Ehrenreich, F. Seitz, D. Turnbull (Academic Press, New York 1972) p. 1.
2. D. R. Hofstadter, Phys. Rev. **B14**, 2239 (1976).
3. D. R. Hofstadter: Gödel, Escher, Bach: an Eternal Golden Braid (Basic books, New York 1979).
4. D. Weiss et al., Surf. Sci. **305**, 408 (1994); R. Schuster et al. in: Proc. 22nd Int. Conf. Phys. Semicond., Vancouver 1994, ed. D. J. Lockwood (World Scientific, Singapore 1995) p. 1979.
5. K. Ensslin, P. M. Petroff, Phys. Rev. **B41**, 12307 (1990); A. Lorke, J. P. Kotthaus, K. Ploog, Phys. Rev. **B44**, 3447 (1991); D. Weiss et al., Phys. Rev. Lett. **66**, 2790 (1991); J. Takahara et al., Jpn. J. Appl. Phys. **30**, 3250 (1991).
6. R. Fleischmann, T. Geisel, R. Ketzmerick, Phys. Rev. Lett. **68**, 1367 (1992).
7. D. Weiss et al., Phys. Rev. Lett. **70**, 4228 (1993).
8. H. Silberbauer, U. Rössler, Phys. Rev. **B50**, 11911 (1994); H. Silberbauer, P. Rotter, M. Suhrke, U. Rössler, Semicond. Sci. Technol. **9**, 1906 (1994).
9. R. Schuster, K. Ensslin in: Festkörperprobleme/Advances in Solid State Physics **34**, ed. by R. Helbig (Vieweg, Braunschweig 1995) p. 195.
10. R. Neudert, P. Rotter, U. Rössler, M. Suhrke, K. Ensslin, J. P. Kotthaus to be published in Phys. Rev. B.
11. K. Richter, Europhys. Lett. **29**, 7 (1995); G. Hackenbroich and F. von Oppen, Europhys. Lett. **29**, 151 (1995).
12. P. Rotter, M. Suhrke, U. Rössler, submitted to Phys. Rev. B.

MAGNETORESISTANCE OF A LATERAL CONTACT TO 2D ELECTRON GAS

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Recently, contacts between superconductors and 2DEG attracted much interest [1, 2, 3, 4, 5]. In this papers, we investigate the $I-V$ characteristics and the contact magnetoresistance in normal state of well-defined *lateral* niobium contacts to high-mobility InGaAs heterostructures (Fig. 1). A strong decrease of the boundary resistance with magnetic field is observed. Such a negative magnetoresistance is known to exist in ballistic constrictions [6], where it is attributed to a reduction of backscattering by the magnetic field. The lateral Nb-2DEG contacts studied in the present paper are no geometrical constrictions in the 2DEG plane but hetero-contacts with a potential barrier located at a sharp interface. We will show that for a hetero-contact an additional mechanism for this negative magnetoresistance is provided by multiple collisions of ballistic electrons with the contact.

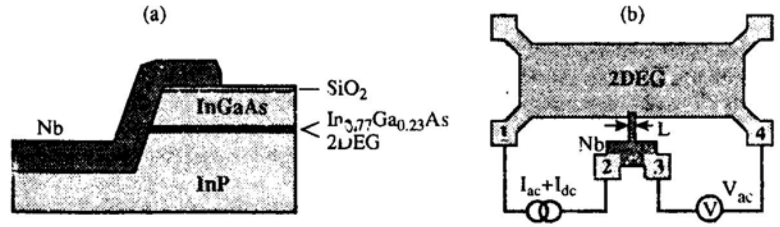

Figure 1: (a) Schematic cross-section of the investigated lateral contact to the 2-dimensional electron gas layer and (b) top view of the sample layout.

Experiments have been performed with InGaAs/InP heterostructures schematically shown in Fig. 1. The 2DEG is formed in a strained $In_{0.77}Ga_{0.23}As$ conduction channel [7, 8]. The In-rich layer allows to obtain a very low Schottky barrier and a vanishing depletion region at the boundary [8, 9]. The conduction channel is placed under a 150 nm thick $In_{0.53}Ga_{0.47}As$ barrier layer. The top surface of the structure has been covered with an insulating layer of SiO_2. Transport measurements yield a mobility of about 3.7×10^5 $cm^2/(V\,s)$ and a sheet carrier concentration of $n_e = 6 \times 10^{11}$ cm^{-2} in the first subband. Lateral contacts to the 10 nm–thick $In_{0.77}Ga_{0.23}As$ layer were prepared by etching the heterostructure and depositing niobium from the side. Prior to the deposition, the surface was cleaned in situ by Ar sputtering.

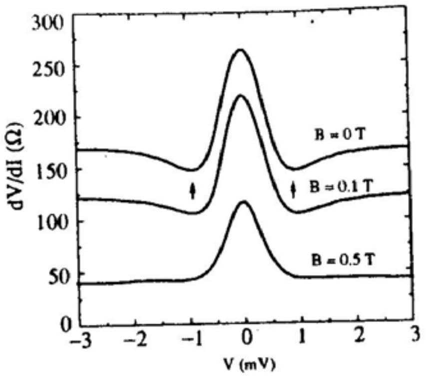

Figure 2: Differential resistance of a 2.4 μm long contact for various magnetic fields B applied perpendicular to the plane of the 2DEG. The measurements have been carried out at 1.2 K with an ac-current of 100 nA.

The differential resistance of the contacts has been measured in the temperature range between 0.05 K and 10 K. A dc-current I_{dc} superimposed by a small ac-current I_{ac} was applied between the contacts 1 and 2 while the ac-voltage drop V_{ac} was detected by lock-in technique between the contacts 3 and 4 [see Fig. 1 (b)]. Contacts 2 and 3 made on superconducting Nb had the same potential resulting in a (quasi) 3-terminal measurement configuration.

A pronounced nonlinearity in the differential resistance $R = dV/dI$ of the contacts due to the superconducting energy gap of Nb has been observed below the transition temperature of Nb (about 9 K). The ratio of the zero-bias resistance R_0 to the normal state resistance R_N of about 1.5 (see Fig. 2) indicates the presence of a well-defined barrier at the lateral interface between the 2DEG and niobium. Increasing the magnetic field from 0 T to 0.5 T leads to a strong decrease of the zero-bias and the normal state resistance of the contact.

The gap voltages V_{gap} of both polarities indicated by arrows in Fig. 2 are smaller than the Nb bulk gap $\Delta_{Nb} \simeq 1.3$ meV due to a thin nonsuperconducting layer which may exist at the Nb/2DEG interface and proximity effect between Nb and 2DEG. To account for these effects a generalization of the BTK model [10] has been done recently.. In the following, we concentrate on the discussion of the normal state properties of our contacts.

In order to determine the boundary resistance, the series resistance of the 2DEG has to be excluded from the data. The used 3-terminal measurement configuration contains a Hall contribution of the 2DEG for one magnetic field direction. For this reason, in our measurements we have chosen the opposite field direction without Hall contribution. Except for the Shubnikov de Haas oscillations at high fields, the 2DEG has a specific sheet resistance of about 25 Ω. Thus, at low fields the series resistance of the 2DEG is definitely smaller than the contact resistance. The zero field resistance might be influenced by an asymmetric current distribution in the 2DEG which is changed in small magnetic fields. We investigated this effect by performing measurements with different contact configurations. The relative change of the resistance for $B = 0$ was in all cases less than 20%. Thus, we conclude that the series resistance of the 2DEG plays a minor role here.

Further, two-dimensional weak localization could be excluded as a possible reason of the negative magnetoresistance: varying the temperature between 50 mK and 10 K had very little influence on the observed $R(B)$ behavior.

Figure 4 shows the experimentally measured resistance (points) of a 0.6 μm long contact as a function of the applied magnetic field B. We observe a strong decrease of the resistance with increasing magnetic field. The measurements have been carried out at a temperature of 50 mK with an ac-current of $I_{ac} = 50$ nA. In order to detect solely the normal state resistance variation, a dc-current of $I_{dc} = 5$ μA has been drawn through the Nb-2DEG contact provided the voltage drop larger than V_{gap}. At fields above 1 T small Shubnikov de Haas oscillations occur which indicate the influence of the 2DEG measured in series with the contact.

All electrons in the 2DEG near the Nb-2DEG interface can be explicitly divided into nonequilibrium current-carrying electrons ("effective electrons" EE) [12, 13, 14] and those in equilibrium which do not contribute to the current. The ballistic mean free path for electron transport in the considered 2DEG is of the order of the contact length L. Therefore, some of the EE move to the contact ballistically while others are scattered in the vicinity of the interface and give rise to a field-independent, diffusive contribution to the current. For the case of a finite transparency of the interface, $\eta < 1$, relevant for our Nb-2DEG contacts, all EE are partly reflected. Depending on the field strength the ballistic EE may return to the contact following cyclotron orbits and give rise to the observed negative magnetoresistance.

The diffusive contribution to the conductance $G(B)$ is given by a modification of the 2D analogue of Sharvin's formula [6, 12, 14]:

$$G_{diff} = M\frac{2e^2}{h}\frac{k_F L}{\pi}\int_0^{\pi/2}\cos(\alpha)D(\alpha)d\alpha \qquad (1)$$

where $M \leq 1$ is the fraction of diffusive EE relative to the total number of initial EE; k_F denotes the Fermi wave vector of the 2DEG electrons and $D(\alpha)$ is the angle dependent transmission coefficient which may be approximated by $D(\alpha) \approx \eta\cos(\alpha)$ for $\eta \ll 1$; η depends on the Fermi velocities and the interface barrier strength; α is the angle between the normal to the boundary and the velocity of the electron (see Fig. 3).

To calculate the ballistic contribution $G_{ball}(B)$ to the conductance we extend the classical analogy to edge state transport namely skipping orbits [15] to very low fields (5 mT). Accordingly the ballistic EE follow cyclotron orbits with a Larmor radius of $r_c = \hbar k_F/(eB)$. We assume further that these electrons are reflected specularly by the sample edge as well as by the contact, as shown in Fig. 3. The angle distribution of the initial ballistic EE is assumed to be homogeneous within the interval $-\pi/2 < \alpha < \pi/2$.

The number of tunneling attempts n for a ballistic EE depends on the ratio $L/(2r_c)$ and directly enters the effective transmission probability $D_{eff}(\alpha, n) = \sum_{i=1}^{n} D(\alpha)(1 - D(\alpha))^{i-1} = 1 - [1 - D(\alpha)]^n$. Therefore, reducing r_c by an increase of the magnetic field leads to an increase of n and consequently of $D_{eff}(\alpha, n)$. The effective transmission probability has to be averaged over angles and starting points of the ballistic trajectories leading to the following expression for the ballistic contribution to $G(B)$:

$$G_{ball}(B) = N\frac{2e^2}{h}\frac{k_F L}{\pi}\int_0^{\pi/2}d\alpha\left\{\cos(\alpha)\frac{1}{s(\alpha)}\int_0^{s(\alpha)}D_{eff}\left(\alpha, \mathrm{Int}\left(\frac{L+x}{s(\alpha)}\right)\right)dx\right\} \qquad (2)$$

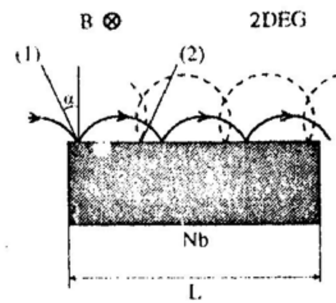

Figure 3: Ballistic electron trajectories at the tunnel barrier interface in the applied magnetic field B for positive (1) and negative (2) angles α.

where $\mathrm{Int}(Y)$ denotes the integer part of Y representing the number of tunneling attempts, and $s(\alpha) = 2r_c\cos(\alpha)$ is the distance between two collisions with the boundary. The constant $N \leq 1$ defines the relative contribution of $G_{\rm diff}$ and $G_{\rm ball}(B)$. As the total number of initial modes is given by $k_F L/\pi$, the normalization condition requires $M + N = 1$.

As follows from Eq. (2) $G_{\rm ball}(B \to 0)$ vanishes, and $G(0)$ can be approximated by $G_{\rm diff}$. Therefore, the normalized magnetoresistance can be written as:

$$\frac{R(B)}{R(0)} = \frac{G(0)}{G(B)} = \frac{G_{\rm diff}}{G_{\rm ball}(B) + G_{\rm diff}} \tag{3}$$

In the considered case of strongly different electronic concentrations of the contacting metals the increased transparency which enters $G_{\rm ball}(B)$ is not compensated by backflow processes from Nb to 2DEG since the Larmor radius in Nb is much larger than the contact length and the mean free path in Nb.

The above discussion of increasing $D_{\rm eff}(\alpha, n)$ by magnetic field in the same way holds for unoccupied states (holes [16]) in the 2DEG: leading to the symmetry of the magnetoresistance with respect to the polarity of the dc-voltage drop V (Fig. 2).

$R(B)/R(0)$ has been calculated numerically from Eqs. (2) and (3) and is shown in Fig. 4. The upper x-axis is directly related to the magnetic field axis taking the geometrical width $L = 0.6$ μm and the known value $B \times r_c = 0.13$ Tμm for our 2DEG. A rather good agreement between experiment and theory can be achieved for the fitting parameters $\eta = 0.1$ and $M/N = 3.4$. The inset shows the corresponding contributions $G_{\rm ball}(B)$ and $G_{\rm diff}$ which add to the total conductance.

For $L/(2r_c) < 1$ the decrease of magnetoresistance is mainly due to the enhancement of the probability for ballistic EE to interact with the contact at all. This effect is comparable to the suppression of backscattering by magnetic field for a geometrical constriction in a 2DEG discussed for the ballistic regime in [6]. For $L/(2r_c) > 1$ all ballistic EE which do not tunnel at their first attempt have more than one tunneling attempt, further decreasing the magnetoresistance. In the

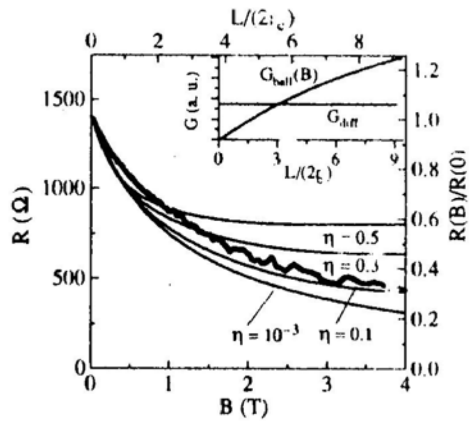

Figure 4: Normal state resistance of a 0.6 μm long contact (points) as a function of the applied magnetic field. The lines show the calculation results using Eq. (3) for values of the barrier transmission coefficient η indicated in the plot. The inset shows $G_{\rm diff}$ and $G_{\rm ball}(B)$ in arbitrary units.

limit of low transparency, $\eta \ll 1$, and for intermediate fields Eq. (2) can be rewritten as

$$G_{\rm ball}(B) = N\frac{2e^2}{h}\frac{k_F L}{\pi}\frac{\eta L}{r_c}, \quad 1 \ll \frac{L}{r_c} \ll \eta^{-1} \tag{4}$$

and $R(B)/R(0)$ becomes independent of η. A corresponding curve for $\eta = 10^{-3}$ is also plotted in Fig. 2. At very high fields when $L/r_c \gg \eta^{-1}$ the magnetoresistance starts to saturate at the level $R(B)/R(0) = M\eta/(M\eta + 2N)$ as $D_{\rm eff}(\alpha, n)$ approaches unity.

We measured more than ten contacts of various dimensions (between 0.6 μm and 100 μm) at different temperatures below and above the critical temperature of Nb. Fair agreement between theory and experiment was found with η values between 0.1 and 0.3.

In summary, we report that Nb contacts to the edge of a high-mobility 2-dimensional electron gas in InGaAs heterostructures show the gap structure of Nb and a large negative magnetoresistance. The latter is explained in terms of a confinement of electron trajectories near the contact boundary with increasing magnetic field. The suggested model shows good agreement with experiment and can be used to characterize the contacts.

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References

[1] T. M. Klapwijk, Physica B 197, 481 (1994).

[2] H. Takayanagi, T. Akazaki and J. Nitta, Phys. Rev. B **51** 1374 (1995).

[3] H. Takayanagi, J. Bindslev Hansen, and J. Nitta, Phys. Rev. Lett. **74**, 162 (1995); Phys. Rev. Lett. **74**, 166 (1995).

[4] H. Kroemer, C. Nguyen and E. L. Hu, Sol.-St. Electr. **37**, 1021 (1994).

[5] K.-M. H. Lenssen, L. A. Westerling, P. C. A. Jeekel, et all, Physica B **194-196**, 2413 (1994).

[6] C. W. J. Beenakker and H. van Houten, Quantum Transport in Semiconductor Nanostructures, In: Solid State Physics, **44**, p.1.

[7] H. Hardtdegen, R. Meyer, H. Løken-Larsen, et all, J. Crystal Growth, **116**, 521, (1992)

[8] H. Hardtdegen, R. Meyer, M. Holfelder, et all, J. Appl. Phys. **73**, 4489 (1993)

[9] J. Appenzeller, Th. Schäpers, H. Hardtdegen et all, Phys. Rev. B **51**, 4336 (1995).

[10] G. E. Blonder, M. Tinkham, and T. M. Klapwijk, Phys. Rev. B **25**, 4515 (1982).

[11] A. A. Golubov and M. Yu. Kupriyanov, Pis'ma Zh. Eksp. Teor. Fiz. **61**, 830 (1995).

[12] Yu. V. Sharvin, Zh. Eksp. Teor. Fiz. **48**, 984 (1965); [JETP **21**, 655 (1965)].

[13] V. S. Tsoi, Pis'ma Zh. Eksp. Teor. Fiz. **19**, 114 (1974); [JETP Lett. **19**, 70 (1974)].

[14] V. F. Gantmakher and I. B. Levinson, Scattering of Carriers in Metals and Semiconductors, Springer, 1986.

[15] M. Büttiker, Phys. Rev. B **38**, 9375 (1988).

[16] J. C. Williamson, H. van Houten, C. W. J. Beenakker et all Surf. Sci. **229**, 303 (1990).

INFLUENCE OF EXCITONIC EFFECTS ON THE TRANSFER OF CHARGE CARRIERS IN ASYMMETRIC DOUBLE QUANTUM WELLS

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Asymmetric double quantum wells (ADQW's) provide an ideal system for studying the transfer of charge carriers through barriers in semiconductor heterostructures. As this transfer (or 'tunneling' in a semi-classical description) is of high interest for fundamental physics as well as for device applications, an abundance of experiments on ADQW has been performed over the last years. Usually, the transfer of electrons and holes is investigated by time resolved photoluminescence excitation experiments: the exciton 1s state formed by the first excited electron and hole subband of the ADQW is resonantly excited by an ultrashort laser pulse. The wave functions belonging to these subbands are located mainly in the narrow well. The time dependence of the luminescence from this exciton as well as the time dependence of the luminescence from the exciton 1s state formed by the lowest electron and hole subband of the ADQW are monitored. From an analysis of these two transients, two time constants can be extracted. They correspond to the transfer rates of electrons and holes. For almost a decade, such experiments have been performed on samples made from III-V-semiconductors. The results could be well explained in the framework of a single particle theory, yielding a fast relaxation rate for the electrons and a rather slow relaxation rate for the holes. The difference in the relaxation rates for the two kinds of charge carriers has its origin in their different effective masses, which in turn leads to a larger subband spacing for the electrons. In many cases, this subband spacing is larger than the energy of a longitudinal optical phonon, while the subband spacing for holes is less than $\hbar\omega_{LO}$. The electrons can therefore relax by the more efficient Fröhlich coupling, while the holes can only relax by the less efficient deformation-potential coupling, i. e. the emission of longitudinal acoustic phonons.

Over the last few years, similar experiments have been performed on samples made from II-VI-semiconductors. In these experiments, an entirely different picture of the tunneling process emerged: electrons and holes do not tunnel independently in a two step process any more (although the respective time constants in the single-particle picture can differ by an order of magnitude). In contrast, the exciton seems to be transfered through the barrier as an entity. The reason for this new behavior has to be found in the larger polarity of the