RAPID COMMUNICATIONS

PHYSICAL REVIEW B VOLUME 46, NUMBER 19 15 NOVEMBER 1992-I

Nonparabolic confinement in quantum wire superlattices

H. Drexler, W. Hansen, and J. P. Kotthaus
Sektion Physik, Universität München, Geschwister-Scholl-Platz 1, 8000 München 22, Germany

M. Holland and S. P. Beaumont
Department of Electronics and Electrical Engineering, University of Glasgow, Glasgow G128QQ, United Kingdom
(Received 8 July 1992)

The high-frequency polarizability of quantum wires in a novel metal-insulator-semiconductor field-effect transistor heterojunction is studied in the far-infrared regime. Whereas on conventional narrow quantum wires only the fundamental mode of the dimensional resonance is observed, the quantum wires investigated here exhibit higher harmonics. Additionally, at high electron densities and intermediate magnetic fields we observe a completely unexpected splitting of the fundamental mode. The behavior is discussed in view of nonparabolic confinement as well as interwire interaction.

Synthetic low-dimensional electron systems with dimensionality below 2 are presently a subject of much interest. For the realization of such systems the two-dimensional (2D) electron gas in metal-oxide-semiconductor systems or heterostructures has turned out to be a very good basis. Much insight into the electronic properties of such systems is gained from investigations of their far-infrared (FIR) excitations. They have demonstrated that in etched,[2] as well as field-effect-induced[3] low-dimensional electron systems, the bare confinement potential V_{ext} is of parabolic shape in a very good approximation. According to the generalized Kohn theorem[4] the FIR conductivity exhibits a single dimensional resonance at a frequency that corresponds to the characteristic frequency of the bare potential. Here we mean by "bare potential" the potential induced by all external charges contributing to the confinement as, e.g., charges on the gate pattern, in surface states or image charges on the semiconductor interfaces, but not the electrons that occupy the quantum wire or dot. Model calculations as well as experiments also demonstrate that the electron systems usually are smaller than the lithographically defined patterns on the surface due to lateral depletion widths of the order of some 10 nm. As a consequence once the dots or wires in a superlattice are defined they can be regarded as decoupled in a very good approximation since wide barriers separate them. On the other hand, for superlattices of strongly interacting dots[5,6] intriguing properties are predicted.

For experimental studies of the effects of nonparabolic terms in the bare potential and of the influence of interaction in a superlattice new types of devices have to be fabricated with a modified confinement scheme. The effect of deviations from parabolic shape in the bare confinement has been observed on Si-MOS (metal-oxide-semiconductor) structures with a very versatile stacked gate configuration.[7] If these structures are operated in the so-called subgrating mode the FIR absorption exhibits several resonances. However, no signatures of interwire interaction are reported and the magnetic-field dispersion has not been investigated.[7] Here we investigate electron channels in a special type of $Al_xGa_{1-x}As/GaAs$ heterojunction induced by positive gate voltages beneath grating stripes which are considerably narrower than in the Si-MOS devices.

Our samples are metal-insulator-semiconductor field-effect transistor (MISFET) AlAs/GaAs heterojunctions.[8] The heterojunction contains a back electrode buried 1.5 μm beneath the heterojunction interface. The gate voltage is applied with respect to this electrode. The gate is formed by a metal grating of period a and an aspect ratio $W/a \approx \frac{2}{3}$, where W is the metal stripe width. At voltages above a threshold voltage of typically $V_{th}=0.9$ V, carriers are injected from the back electrode into the potential minima at the heterojunction interface. The 32 nm thick barrier is formed by an undoped AlAs/GaAs superlattice and growth is finished by 10 nm undoped GaAs so that at $V_g > V_{th}$ we expect a capacitance of roughly 1.3×10^{12} e/cm^2 V beneath a homogeneous gate.[8]

Typical FIR transmission spectra on a device with a grating gate of $a=350$ nm are sketched in Fig. 1. The radiation is incident normal onto the sample surface and polarized normal to the grating stripes. At a frequency of about $\bar{v}_1=30$ cm^{-1} a strong resonance is observed that hardly changes position with gate voltage. This resonance is the fundamental mode and is the only excitation observed for electron systems bound in a purely parabolic bare potential. It corresponds to the center-of-mass motion of all confined electrons in a homogeneous electric field. In presence of nonparabolic terms to the confinement a redistribution of the FIR oscillator strengths is expected so that besides the fundamental mode higher harmonics of the dimensional resonance become visible.[9] Indeed in the spectra of Fig. 1 a weak but distinct resonance is found at higher frequencies of about 60–70 cm^{-1}. The oscillator strength of this mode normalized to the corresponding oscillator strength of the fundamental mode first increases at low gate voltages above threshold and then decreases again with a maximum at gate bias $V_g - V_{th}=0.25$ V. From this we infer that the nonparabolicity of the bare potential is gate-

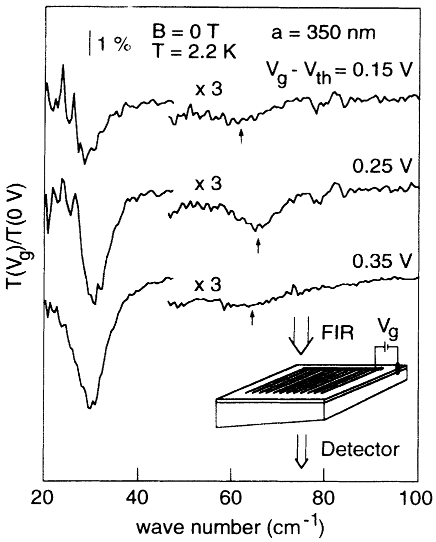

FIG. 1. FIR transmission spectra recorded on a heterojunction MISFET with a grating gate of $a = 350$ nm period and gate biases of 0.15, 0.25, and 0.35 V above the inversion threshold. The experimental setup is sketched in the inset.

voltage dependent.

The resonance frequency of the nth-order dimensional resonance can be expected at[8,10]

$$\frac{\bar{\nu}_n}{\bar{\nu}_1} = \left[n \frac{1+\coth(q_1 d)}{1+\coth(n q_1 d)} \right]^{1/2}, \quad n = 1,2,3,\ldots .$$

Here n is the order of the mode, $\bar{\nu}_1$ is the frequency of the fundamental mode, $q_1 = 2\pi/a$, and d is the separation between the gate and the electron system. The second factor of the right-hand side reflects the screening of the charge-density excitations by the gate, which is less effective at higher wave vectors $\mathbf{n}\cdot\mathbf{q}_1$ that can be approximately assigned to the nth-order mode.[11] Even modes are optically inactive[9] and the $n = 3$ mode is expected within this model at a frequency $\bar{\nu}_3 = 1.95\bar{\nu}_1$ above the fundamental mode. This is close to the experimental values.

We have investigated the dimensional resonances in samples with different periods a of the wire array. In samples with period $a = 500$ nm the resonance frequency of the fundamental mode is slightly (10%) lower compared to the resonances in the sample with $a = 350$ nm. The relative oscillator strength of the higher-order mode is larger in the 500 nm samples. Since our different samples have approximately the same aspect ratio W/a the latter indicates that the nonparabolic contributions become smaller with decreasing widths of the gate stripes. In samples with $a = 250$ nm the resonance frequencies of the fundamental mode are considerably larger: $\bar{\nu}_1 = 50$ cm^{-1}. Imperfections in this sample cause low oscillator strengths so that a higher harmonic mode has not been identified. In general, the relative oscillator strengths of the higher-order modes are considerably weaker in our samples than in wide electron wires in stacked gate Si-MOS channels.[7] It is thus tempting to conclude that the deviations from parabolic bare potential are small in our electron wires. If the bare potential is parabolic an analytical ansatz for the charge-density distribution can be used for an estimate of the effective wire width and one-dimensional charge density N_L from the experimental resonance positions. The model[10] assumes a classical 2D charge-density distribution. Input parameters are the maximum value of the density distribution and an effective dielectric constant $\bar{\epsilon}$ of the background. With $\bar{\epsilon} = 16.4$ and a maximum density calculated from the capacitance value given above we get, e.g., from the experimental resonance position in the samples with period $a = 350$ nm a wire width of $W_{\mathrm{eff}} = 200$ nm and $N_L = 7\times 10^6$ cm^{-1} at $V_g - V_{\mathrm{th}} = 0.35$ V.

In Fig. 2 we depict the resonance positions measured on the same device as function of a magnetic field B applied perpendicular to the sample surface. The dispersion of $\bar{\nu}_1(B)$ is the same as measured previously in different wire arrays:[3] $\bar{\nu}_1(B)^2 = \bar{\nu}_1(0)^2 + \bar{\nu}_c^2$. Here $\bar{\nu}_c$ is the cyclotron-resonance frequency. It is remarkable that the fundamental mode is not affected at all when it crosses the second harmonic of the cyclotron resonance. In contrast the $n = 3$ mode exhibits a very pronounced anticrossing with the second harmonic of the cyclotron resonance. We attribute this behavior to the different characters of the two modes. The fundamental mode originates from the center-of-mass motion of the electron system and is not associated with internal charge oscillations. This uniform mode is not expected to interact with the cyclotron mode.[12] In contrast the higher-order modes are associated with internal oscillations of the many-particle system. In this sense the higher-order modes are analogous to charge-density excitations in 2D systems which are known to interact with the second harmonic of the cyclotron resonance.[12,13] Similar behavior has also been found for one-dimensional plasmons, i.e., charge-density excitations with the wave vector along narrow electron wires.[14]

We observe at high electron densities a completely unexpected behavior of the fundamental mode. In the sample of Fig. 2 the fundamental mode behaves only as discussed above at gate voltages below $V_g - V_{\mathrm{th}} = 0.35$ V.

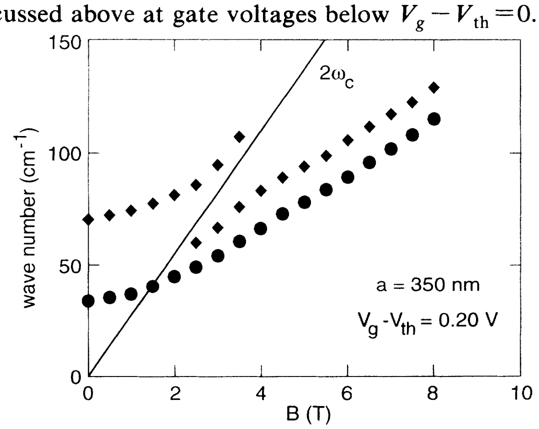

FIG. 2. Positions of the dimensional resonances measured as function of a magnetic field applied perpendicular to the sample surface. Circles denote the positions of the fundamental; rhombs the positions of the $n = 3$ mode. Positions calculated for the second harmonic of the cyclotron resonance are indicated by the straight line.

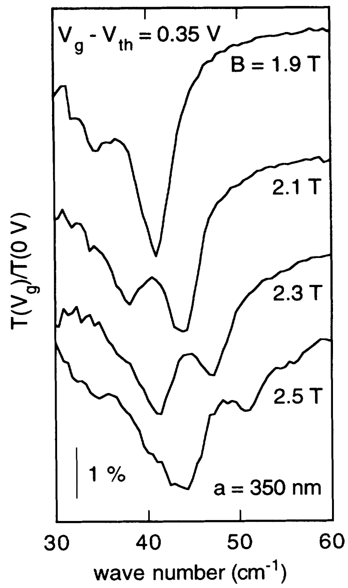

FIG. 3. FIR transmission spectra of the fundamental of the dimensional resonance recorded on the sample of Fig. 2 at higher electron density.

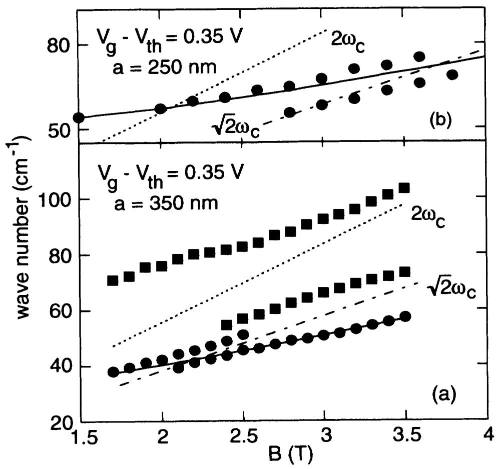

FIG. 4. Resonance positions vs magnetic field recorded in the magnetic-field range where the fundamental of the dimensional resonance splits on (*a*) the sample of Fig. 3 and (*b*) a sample with period $a=250$ nm. Circles denote the positions of the fundamental; rhombs those of the higher-order mode. Positions calculated for the second harmonic of the cyclotron resonance are denoted by dashed lines. Full lines are calculated according to $\bar{\nu}_1(B)^2=\bar{\nu}_1^2+\bar{\nu}_c^2$ with frequencies $\bar{\nu}_1=30$ and 50 cm^{-1} extracted from the experiment in (*a*) and (*b*), respectively. Dashed-dotted straight lines are calculated according to $\bar{\nu}(B)=\sqrt{2}\bar{\nu}_c$.

At higher gate bias the fundamental mode splits in a small magnetic-field range and shows an anticrossing behavior of the resonance oscillator strengths as demonstrated in Fig. 3. The resonance positions in the corresponding magnetic-field regime of samples with periods $a=250$ and $a=350$ nm are plotted in Fig. 4. The strength of the splitting and the magnetic-field range are independent of the polarization of the radiation with respect to the wires. As indicated by the dashed lines the resonance positions of the vanishing branches in both samples converge to $\bar{\nu}=\sqrt{2}\bar{\nu}_c$ suggesting an anticrossing with a mode of corresponding magnetic-field dispersion. However, so far no such excitation has been predicted for an electron wire array. Apart from the range where the fundamental mode splits the magnetic-field dispersion is still well described by $\bar{\nu}_1(B)^2=\bar{\nu}_c^2+\bar{\nu}_1(0)^2$ as indicated by the full line in Fig. 4. The unexpected behavior is only observed if the gate bias exceeds a critical value that decreases with decreasing array period. In samples with $a=500$ nm the splitting occurs at voltages $V_g-V_{th}\geq 0.4$ V. In devices with $a=250$ nm at voltages below $V_g-V_{th}=0.3$ V the signal-to-noise ratio of the spectra does not allow us to identify the presence of a splitting. At higher gate voltages the oscillator strength is sufficiently large to detect a splitting.

At present we do not have a model that describes the splitting of the fundamental mode. From the presence of the $n=3$ mode of the dimensional resonance one could infer that the nonparabolicity of the bare confinement also was responsible for the splitting. One thus might invoke an anomaly of the fundamental mode once the extension of the classical motion perpendicular to the wire axis in a pure harmonic potential $2\hbar k_F/m^*(\omega_0^2+\omega_c^2)^{1/2}$ becomes smaller than a critical width determined by the strength of the nonparabolic contributions to the actual bare potential. Power spectra of numerically evaluated classical orbits[15] in a model potential consisting of a harmonic potential truncated by infinitely steep walls actually exhibit two distinct frequencies that behave similar to our experimental data. This indicates that nonparabolic confinement may be the cause of the splitting although a sophisticated calculation of the magnetic-field dispersion has to include intrawire electron-electron interaction even for the fundamental mode. We note that with the potential models used so far for such calculations no such splitting of the fundamental mode neither in semiclassical[9] nor in quantum-mechanical[16] theories are predicted.

We also consider interwire interaction as a possible origin of the splitting. Since interwire Coulomb interaction increases with decreasing wire separation and increasing electron density[10] this would be in correspondence with our observations. However, it would imply a sudden change of the interaction at a critical magnetic field and we do not understand why the magnetic field should have such a dramatic effect on the interaction potential. Calculations of collective excitations in an array of interacting wires published so far do not consider a magnetic field.[17,18]

Note added in proof. We have learned recently that a similar splitting of the fundamental mode in a magnetic field has also been observed by Gerhardts *et al.*[19] in different samples.

We gratefully acknowledge helpful discussions with C. Dahl and financial support by the Deutsche Forschungsgemeinschaft and the Esprit Basic Research Action.

[1] For a recent review see *Semiconductors and Semimetals,* edited by R. K. Williardson, A. C. Beer, and E. R. Weber (Academic, San Diego, 1992) Vol. 35.

[2] T. Demel, D. Heitmann, P. Grambow, and K. Ploog, Superlatt. Microstruct. **9**, 285 (1991).

[3] W. Hansen, in *Quantum Coherence in Mesoscopic Systems,* Vol. 254 of *NATO Advanced Studies Institute, Series B: Physics,* edited by B. Kramer (Plenum, New York, 1991), p. 23.

[4] S. K. Yip, Phys. Rev. B **43**, 1707 (1991), and references therein.

[5] P. Bakshi, D. A. Broido, and K. Kempa, J. Appl. Phys. **70**, 5150 (1991).

[6] A. V. Chaplik and L. Ioriatti, Surf. Sci. **263**, 354 (1992).

[7] J. Alsmeier, E. Batke, and J. P. Kotthaus, Phys. Rev. B **40**, 12 574 (1989).

[8] H. Drexler, W. Hansen, J. P. Kotthaus, M. Holland, and S. P. Beaumont, Semicond. Sci. Technol. **7**, 1008 (1992).

[9] G. Eliasson, J.-W. Wu, P. Hawrylak, and J. J. Quinn, Solid State Commun. **60**, 41 (1986).

[10] V. B. Shikin, T. Demel, and D. Heitmann, Zh. Eksp. Teor. Fiz. **96**, 1406 (1989) [Sov. Phys. JETP, **69**, 797 (1989)].

[11] A. V. Chaplik, Zh. Eksp. Teor. Fiz. **62**, 746 (1972) [Sov. Phys. JETP **35**, 395 (1972)].

[12] A. V. Chaplik and D. Heitmann, J. Phys. C**18**, 3357 (1985).

[13] E. Batke, D. Heitmann, J. P. Kotthaus, and K. Ploog, Phys. Rev. Lett. **54**, 2367 (1985).

[14] T. Demel, D. Heitmann, P. Grambow, and K. Ploog, Phys. Rev. Lett. **66**, 2657 (1991).

[15] A. Lorke, Surf. Sci. **263**, 307 (1992).

[16] Q. P. Li and S. Das Sarma, Phys. Rev. B **44**, 6277 (1991).

[17] Q. P. Li and S. Das Sarma, Phys. Rev. B **43**, 11 768 (1991), and references therein.

[18] A. Gold, Philos. Mag. Lett. **66**, 163 (1992).

[19] R. R. Gerhardts, P. Grambow, V. Gudmundsson, D. Heitmann, B. Meurer, D. Pfannkuche, and K. Ploog (unpublished).