QUANTUM COMPUTING AND QUANTUM COMMUNICATION WITH ELECTRONS IN NANOSTRUCTURES

Daniel Loss, Guido Burkard and Eugene V. Sukhorukov

Department of Physics and Astronomy, University of Basel,
Klingelbergstrasse 82, CH-4056 Basel, Switzerland

If the states of spins in solids can be created, manipulated, and measured at the single-quantum level, an entirely new form of information processing, quantum computing and quantum communication, will be possible. We review our proposed spin-quantum dot architecture for a quantum computer and review some recent results on a deterministic source of entanglement generated by coupling quantum dots. Addressing the feasibility of quantum communication with entangled electrons we consider a scattering set-up with an entangler and beam splitter where the current noise exhibits bunching behavior for electronic singlet states and antibunching behavior for triplet states. We show that spin currents can produce noise even in the absence of any charge currents.

1 Introduction

Quantum computation (QC) has attracted much interest recently, as it opens up the possibility of outperforming classical computation through new and more powerful quantum algorithms such as the ones discovered by Shor and by Grover. There is now a growing list of quantum tasks such as secret sharing, error correction schemes, quantum teleportation, etc., that have indicated even more the desirability of experimental implementations of QC. In QC the state of each bit is allowed to be any state of a quantum two-level system—a qubit, and QC proceeds by one and two-qubit gates by which all quantum algorithms can be implemented. There are now a number of schemes which have been proposed to realize physical implementations of qubits and quantum gates. Here we will review our qubit proposal based on the spin of electrons confined in quantum dots (or other confined structures such as molecules or atoms). One of the essential features of such qubits is that they are scalable to many qubits, that recent experiments demonstrated very long spin decoherence times in semiconductors, and that the qubit defined as electron-spin is mobile and thus can be used for implementing quantum communication schemes.

We first review our recent results on quantum gates and quantum dots by which entanglement can be generated, and then we describe a set-up by which such entanglement can be observed in transport/noise measurements.

2 Quantum Computing: Electron Spin as Qubit in Quantum Dots

Semiconductor quantum dots, sometimes referred to as artificial atoms, are small devices in which charge carriers are confined in all three dimensions. The confinement is usually achieved by electrical gating and/or etching techniques applied e.g. to a two-dimensional electron gas (2DEG). In GaAs heterostructures the number of electrons in the quantum dots can be changed one by one starting from zero. Typical magnetic fields \( B \approx 1 \text{T} \) correspond to magnetic lengths on the order of \( l_B \approx 10 \text{nm} \), being of the same size as quantum dots. As a consequence, the dot spectrum depends strongly on the applied magnetic field. In coupled quantum dots which can be considered as artificial molecules, Coulomb blockade effects and magnetization have been observed as well as the formation of a delocalized “molecular state”.

In addition to a well defined qubit such as the electron spin considered here, we also need a deterministic “source of entanglement”, i.e. a mechanism by which two specified qubits at a time can be entangled so as to produce the fundamental quantum XOR (or controlled-NOT)
gate operation, represented by a unitary operator $U_{\text{XOR}}$. This can be achieved by temporarily coupling two spins. In short, due to the Coulomb interaction and the Pauli exclusion principle the ground state of two coupled electrons is a spin singlet, i.e. a highly entangled spin state. This physical picture translates into an exchange coupling $J(t)$ between the two spins $S_1$ and $S_2$ described by a Heisenberg Hamiltonian

$$H_s(t) = J(t) \cdot S_1 \cdot S_2.$$  

If the exchange coupling is pulsed such that $\int dtJ(t)/\hbar = J_0\tau_s/\hbar = \pi \mod 2\pi$, the associated unitary time evolution $U(t) = T \exp(i \int_0^t H_s(\tau)d\tau/\hbar)$ corresponds to the “swap” operator $U_{\text{sw}}$ which simply exchanges the quantum states of qubit 1 and 2. Furthermore, the quantum XOR can be obtained by applying the sequence $\exp(i(\pi/2)S_1^z) \exp(-i(\pi/2)S_2^z)U_{\text{sw}}^{1/2} \exp(i\pi S_1^z)U_{\text{sw}}^{1/2} \equiv U_{\text{XOR}}$, i.e. a combination of “square-root of swap” $U_{\text{sw}}^{1/2}$ and single-qubit rotations $\exp(i\pi S_1^z)$, etc. Since $U_{\text{XOR}}$ (combined with single-qubit rotations) is proven to be a universal quantum gate, it can therefore be used to assemble any quantum algorithm. Thus, the study of a quantum XOR gate is essentially reduced to the study of the exchange mechanism and how the exchange coupling $J(t)$ can be controlled experimentally. We wish to emphasize that the switchable coupling mechanism described in the following need not be confined to quantum dots: the same principle can be applied to other systems, e.g. coupled atoms in a Bravais lattice, supramolecular structures, or overlapping shallow donors in semiconductors. As to the latter example, it has been demonstrated in Ref. that the coupling mechanism proposed by us can also be used to couple nuclear spins (playing the role of qubits) via a tunable exchange coupling of the $s$-electrons (of P embedded in Si) to which the nuclear spins are coupled via hyperfine interaction.

The main reason to concentrate here on quantum dots is that these systems are at the center of many ongoing experimental investigations in mesoscopic physics, and thus there seems to be reasonable hope that these systems can be made into quantum gates function along the lines described here.

### 2.1 Model for the Quantum Gate with Coupled Quantum Dots

We consider a system of two laterally coupled quantum dots containing one (conduction band) electron each. It is essential that the electrons are allowed to tunnel between the dots, and that the total wave function of the coupled system must be antisymmetric. It is this fact which introduces correlations between the spins via the charge (orbital) degrees of freedom. The Hamiltonian for the coupled system is $H = \sum_{i=1,2} h_i + C + H_Z = H_{\text{orb}} + H_Z$, where

$$h_i = \frac{1}{2m} \left( p_i - \frac{e}{\kappa} A(r_i) \right)^2 + V(r_i), \quad C = \frac{e^2}{\kappa |r_1 - r_2|}. \quad (2)$$

Here, $h_i$ describes the single-electron dynamics confined to the $xy$-plane, with $m = 0.067 m_e$ being the effective mass in GaAs and $S_i$ the electron spin. The dielectric constant in GaAs is $\kappa = 13.1$. We allow for a magnetic field $B = (0,0,B)$ applied along the $z$-axis and which couples to the electron charge via the vector potential $A(r) = \frac{B}{2} (-y,x,0)$, and to the spin via a Zeeman coupling term $H_Z$. The coupling of the dots (which includes tunneling) is modeled by a quartic potential, $V(x,y) = \frac{m_0}{2} \left( \frac{1}{4\pi^2} (x^2 - a^2)^2 + y^2 \right)$, which separates into two harmonic wells of frequency $\omega_0$, one for each dot, in the limit $2a \gg 2a_B$, where $a$ is half the distance between the centers of the dots, and $a_B = \sqrt{\hbar/m \omega_0}$ is the effective Bohr radius of a dot. This choice for the potential is motivated by the experimental fact that the spectrum of single dots in GaAs is well described by a parabolic confinement potential, e.g. with $\hbar \omega_0 = 3 \text{meV}$, $a_B = 20 \text{nm}$.

The (bare) Coulomb interaction between the two electrons is described by $C$. The screening length $\lambda$ in almost depleted regions like few-electron quantum dots can be expected to be much
larger than the bulk 2DEG screening length (which is about 40 nm in GaAs). Therefore, $\lambda$ is large compared to the size of the coupled system, $\lambda \gg 2a \approx 40$ nm for small dots, and we will consider the limit of unscreened Coulomb interaction ($\lambda/a \gg 1$).

### 2.2 Exchange Interaction in the Heitler-London approach

At low-temperatures where $kT \ll \hbar \omega_0$ we can restrict ourselves to the two lowest orbital eigenstates of $H_{\text{orb}}$, one of which is symmetric (spin singlet) and the other one antisymmetric (spin triplet). In this reduced (four-dimensional) Hilbert space, $H_{\text{orb}}$ can be replaced by the effective Heisenberg spin Hamiltonian Eq. (2), where the exchange energy $J = \epsilon_t - \epsilon_s$ is the difference between the triplet and singlet energy. The above model cannot be solved in an analytically closed form. However, the analogy between atoms and quantum dots (artificial atoms) provides us with a powerful set of variational methods from molecular physics for finding $\epsilon_t$ and $\epsilon_s$. In the Heitler-London approximation and making use of the Darwin-Fock solution for the isolated dots we find,

$$J = \frac{\hbar \omega_0}{\sinh(2d^2(2b - 1/b))} \left[ c\sqrt{b} \left( e^{-bd^2}I_0(bd^2) - e^{d^2(b-1/b)}I_0(d^2(b-1/b)) \right) + \frac{3}{4b} \left( 1 + bd^2 \right) \right], \quad (3)$$

where we introduce the dimensionless distance $d = a/a_B$, and $I_0$ is the zeroth order Bessel function. The first and second terms in Eq. (3) are due to the Coulomb interaction $C$, where the exchange term enters with a minus sign. The parameter $c = \sqrt{\pi/2(e^2/\kappa a_B)}/\hbar \omega_0$ ($\approx 2.4$, for $\hbar \omega_0 = 3$ meV) is the ratio between Coulomb and confinement energy. The last term comes from the confinement potential $W$. Note that typically $|J/\hbar \omega_0| \lesssim 0.2$. Also, we see that $J > 0$ for $B = 0$, which must be the case for a two-particle system that is time-reversal invariant. The most remarkable feature of $J(B)$, however, is the change of sign from positive to negative at $B = B_s$ (which occurs over a wide range of parameters $c$ and $a$). This singlet-triplet crossing occurs at about $B_s = 1.3$ T, for $\hbar \omega_0 = 3$ meV ($c = 2.42$) and $d = 0.7$. The transition from antiferromagnetic ($J > 0$) to ferromagnetic ($J < 0$) spin-spin coupling with increasing magnetic field is caused by the long-range Coulomb interaction, in particular by the negative exchange term, the second term in Eq. (3). As $B \gg B_0$ ($\approx 3.5$ T for $\hbar \omega_0 = 3$ meV), the magnetic field compresses the orbits by a factor $b \approx B/B_0 \gg 1$ and thereby reduces the overlap of the wavefunctions exponentially. Similarly, the overlap decays exponentially for $d \gg 1$. Note however, that this exponential suppression is partly compensated by the exponentially growing exchange term $\propto \exp(2d^2(b-1/b))$. As a result, $J$ decays exponentially as $\exp(-2d^2b)$ for large $b$ or $d$. Thus, $J$ can be tuned through zero and then exponentially suppressed to zero by a magnetic field in a very efficient way (exponential switching is highly desirable to minimize gate errors). This sign reversal of $J$ is due to the long-range Coulomb forces and is not contained in the standard Hubbard approximation which takes only short-range interactions into account. We note that our Heitler-London approximation breaks down explicitly (i.e. $J$ becomes negative even when $B = 0$) for certain inter-dot distances when $c$ exceeds 2.8. By working around the magnetic field where $J$ vanishes the exchange interaction can be pulsed on, even without changing the tunneling barrier between the dots, either by an application of a local magnetic field, or by exploiting a Stark electric field (which will also make the exchange interaction nonzero).

Qualitatively similar results are obtained when we refine above Heitler-London result by taking into account higher levels and double occupancy of the dots (requiring a molecular orbit approach), and also for vertically coupled dots. Finally, we note that a spin coupling can also be achieved on a long distance scale by using a cavity-QED scheme.
3 Quantum Communication with Electrons: Detection of Entanglement

The availability of pairwise entangled qubits—Einstein-Podolsky-Rosen (EPR) pairs—is a necessary prerequisite in quantum communication. The prime example of an EPR pair considered here is the singlet state formed by two electron spins, its main feature being its non-locality: If we separate the two electrons in real space, their total spin state can still remain entangled. Such non-locality gives rise to striking phenomena such as violations of Bell inequalities and quantum teleportation and has been investigated for photons, but not yet for massive particles such as electrons, let alone in a solid state environment. This is so because it is difficult to first produce and to second detect entanglement of electrons in a controlled way. In this section we describe an experimental set-up by which the entanglement of electrons (once produced as described in the previous section) can be detected in noise measurements, see Fig. 1. The entangler is assumed to be a device by which we can generate (or detect) entangled electron states, a specific realization being above-mentioned double-dot system. The presence of a beam splitter ensures that the electrons leaving the entangler have a finite amplitude to be interchanged (without mutual interaction). The quantity of interest is then the current-current correlations (noise) measured in leads 3 and/or 4. It is well-known that particles with symmetric wave functions show bunching behavior in the noise, whereas particles with antisymmetric wave functions show antibunching behavior. The latter situation is the one considered so far for electrons in the normal state both in theory and in experiments. However, since the noise is produced by the charge degrees of freedom we can expect that in the absence of spin scattering processes the noise is sensitive to the symmetry of only the orbital part of the wave function. On the other hand, since the spin singlet of two electrons is uniquely associated with a symmetric orbital part

\[ I_{\alpha\sigma}(t) = \frac{e}{\hbar \nu} \sum_{E,E'} \left[ a_{\alpha\sigma}^\dagger(E) a_{\alpha\sigma}(E') - b_{\alpha\sigma}^\dagger(E) b_{\alpha\sigma}(E') \right] \exp\left[ i(E - E') t / \hbar \right], \]

where \( a_{\alpha\sigma}^\dagger(E) \) creates an incoming electron in lead \( \alpha \) with spin \( \sigma \) and energy \( E \), and the operators \( b_{\alpha\sigma} \) for the outgoing electrons are related to the operators \( a_{\alpha\sigma} \) for the incident electrons via the scattering matrix, \( s_{\alpha\beta}, b_{\alpha\sigma}(E) = \sum_{\beta} s_{\alpha\beta} a_{\beta\sigma}(E) \). We will assume that the scattering matrix is spin- and energy-independent. Note that since we are dealing with discrete energy states here, we normalize the operators \( a_{\alpha\sigma}(E) \) such that \( [a_{\alpha\sigma}(E), a_{\beta\sigma'}(E')]^\dagger = \delta_{\sigma\sigma'} \delta_{\alpha\beta} \delta_{E,E'}/\nu \), where \( \delta_{E,E'} \) is the Kronecker symbol, and \( \nu \) the density of states. We assume that each lead consists of only a single quantum channel. We then obtain for the current

\[ I_{\alpha\sigma}(t) = \frac{e}{\hbar \nu} \sum_{E,E'} \sum_{\beta,\gamma} a_{\beta\sigma}^\dagger(E) A_{\beta\gamma} a_{\gamma\sigma}(E') e^{i(E - E') t / \hbar}, \quad A_{\beta\gamma} = \delta_{\alpha\beta} \delta_{\alpha\gamma} - s_{\alpha\beta}^* s_{\alpha\gamma}. \]

We restrict ourselves here to unpolarized currents, \( I_{\alpha} = \sum_{\sigma} I_{\alpha\sigma} \). The spectral density of the current fluctuations (noise) \( \delta I_{\alpha} = I_{\alpha} - \langle I_{\alpha} \rangle \) between the leads \( \alpha \) and \( \beta \) are defined as

\[ S_{\alpha\beta}(\omega) = \lim_{T \to \infty} \frac{h \nu}{T} \int_0^T dt \ e^{i\omega t} \langle \delta I_{\alpha}(t) \delta I_{\beta}(0) \rangle |\Psi\rangle, \]

where the state \( |\Psi\rangle \) is some arbitrary state to be specified below. Inserting the expression for...
the currents Eq. (7) into this definition, we obtain for the zero frequency correlations

\[ S_{\alpha\beta} = \frac{e^2}{h \nu} \sum_{\gamma\delta\zeta} A^\alpha_{\gamma\delta} A^\beta_{\zeta} \sum_{E, E', \sigma\sigma'} \left[ \langle \Psi | a^\dagger_{\gamma\sigma}(E)a_{\delta\sigma}(E) a^\dagger_{\epsilon\sigma'}(E') a_{\zeta\sigma'}(E') | \Psi \rangle - \langle \Psi | a^\dagger_{\epsilon\sigma'}(E') a_{\zeta\sigma'}(E') | \Psi \rangle \langle \Psi | a^\dagger_{\gamma\sigma}(E)a_{\delta\sigma}(E) | \Psi \rangle \right] . \]

(7)

We note that since |Ψ⟩ in general does not describe a Fermi liquid state, it is not possible to apply Wick’s theorem.

We will now investigate the noise correlations for scattering with the entangled incident state |Ψ⟩ = |±⟩, where |±⟩ = 1/\sqrt{2} \left( a^\dagger_2(e_2)a_1(e_1) ± a^\dagger_1(e_2)a_1^\dagger(e_1) \right) |0⟩ . The state |−⟩ is the spin singlet, |S⟩, while |+⟩ denotes one of the spin triplets |T_{0,±}⟩; in the following we will present a calculation of the noise for |+⟩ = |T_0⟩, i.e. the triplet with m_z = 0.

Substituting |±⟩ for |Ψ⟩, we get \langle ± | \delta I_\alpha \delta I_\beta | ± \rangle = \langle ↑↓ | \delta I_\alpha \delta I_\beta | ↑↓ \rangle ± \langle ↑↓ | \delta I_\alpha \delta I_\beta | ↓↑ \rangle , where the upper (lower) sign of the exchange term refers to triplet (singlet). After some straightforward manipulations, we obtain the following result for the correlations between the leads \alpha and \beta, \( S_{\alpha\beta} = \frac{e^2}{2h\nu} \sum_{\gamma\delta\zeta} A^\alpha_{\gamma\delta} A^\beta_{\zeta} \delta_{\epsilon_1,\epsilon_2} \left( A^\dagger_{12} A^\dagger_{12} + A^\dagger_{21} A^\dagger_{12} \right) \), where \( \gamma, \delta \) denotes the sum over \( \gamma = 1, 2 \) and all \( \delta \neq \gamma \), and where again the upper (lower) sign refers to triplets (singlets).

We apply above formula now to the set-up shown in Fig. 1 involving four leads, described by the scattering matrix elements, s_{31} = s_{42} = r, and s_{41} = s_{32} = t, where r and t denote the reflection and transmission amplitudes at the beam splitter, resp., and with no backscattering, s_{12} = s_{34} = s_{\alpha\alpha} = 0. The unitarity of the s-matrix implies |r|^2 + |t|^2 = 1, and Re[r^*t] = 0. Using above relations, we obtain finally:

\[ S_{33} = S_{44} = -S_{34} = 2 \frac{e^2}{h \nu} T (1 - T) (1 ± \delta_{\epsilon_1,\epsilon_2}) , \]

(8)

where T = |t|^2 is the probability for transmission through the beam splitter. The calculation for the remaining two triplet states |+⟩ = |T_{0,±}⟩ = |↑↓⟩, |↓↑⟩ yields the same result Eq. (8) (upper sign). Note that the total current \( \delta I_3 + \delta I_4 \) does not fluctuate, i.e. \( S_{33} + S_{44} + 2S_{34} = 0 \), since we have excluded backscattering. For the average current in lead \alpha we obtain \( \langle I_{\alpha} \rangle = e/\hbar \nu \), with no difference between singlets and triplets. Then, the Fano factor \( F = S_{\alpha\alpha}/|\langle I_{\alpha} \rangle| \) takes the following form

\[ F = 2eT(1 - T) (1 ± \delta_{\epsilon_1,\epsilon_2}) , \]

(9)

and correspondingly for the cross correlations. This result confirms our expectation stated in the introduction: if two electrons with the same energies, \( \epsilon_1 = \epsilon_2 \), in the singlet state |S⟩ = |−⟩ are injected into the leads 1 and 2, then the zero frequency noise is enhanced by a factor of two, \( F = 4eT(1 - T) \), compared to the shot noise of uncorrelated particles, \( F = 2eT(1 - T) \). This enhancement of noise is due to bunching of electrons in the outgoing leads, caused by the symmetric orbital wavefunction of the spin singlet |S⟩. On the other hand, the triplet states |+⟩ = |T_{0,±}⟩ exhibit an antibunching effect, leading to a complete suppression of the zero-frequency noise in Eq. (8), \( S_{\alpha\alpha} = 0 \). The noise enhancement for the singlet |S⟩ is a unique signature for entanglement (there exists no unentangled state with the same symmetry), therefore entanglement can be observed by measuring the noise power of a mesoscopic conductor. The triplets can be further distinguished from each other if we can measure the spin of the two electrons in the outgoing leads, or if we insert spin-selective tunneling devices into leads 3,4 which would filter a certain spin polarization.

We emphasize that above results remain unchanged if we consider states |±⟩ which are created above a Fermi sea. We have shown elsewhere that the entanglement of two electrons propagating in a Fermi sea gets reduced by the quasiparticle weight z_F (for each lead one factor) due to the presence of interacting electrons. In the metallic regime z_F assumes typically some
Figure 1: Uncorrelated electrons are fed into the entangler $E$ through the Fermi leads $1'$ and $2'$. The entangler is a device (see text) that produces pairs of electrons in the entangled spin singlet or one of the spin triplets and injects one of the electrons into lead 1 and the other into lead 2. For singlets we get bunching due to their orbital symmetry, whereas for triplets we get antibunching due to their orbital antisymmetry. Note that instead of being divided into an entangler and a beam splitter, the whole area inside the dotted box can be viewed as a two-particle scattering region.

Finally we discuss the current noise induced by the spin transport in a two-terminal conductor attached to Fermi leads with spin-dependent chemical potentials $\mu_{\sigma}$. From Eq. (5) we immediately obtain $\langle I_{\sigma} \rangle = eT \Delta_{\mu_{\sigma}}$, where we have introduced the difference of chemical potentials $\Delta_{\mu_{\sigma}} = \mu_{1\sigma} - \mu_{2\sigma}$ for each spin orientation $\sigma$. Then, applying Wick’s theorem to Eq. (7) we obtain the noise power $S = \frac{e^2}{h} T (1 - T) (|\Delta_{\mu_{\uparrow}}| + |\Delta_{\mu_{\downarrow}}|)$, thus the contribution of two spin subsystems to the noise is independent, as it should be if there is no spin interaction. Therefore, we can rewrite above expression as $S = e(1 - T) (|\langle I_{\uparrow} \rangle| + |\langle I_{\downarrow} \rangle|)$. In particular, when $\Delta_{\mu_{\uparrow}} = \Delta_{\mu_{\downarrow}}$ we obtain the usual result $S = e(1 - T) |I_c|$ for the shot noise induced by the charge current $I_c \equiv \langle I_{\uparrow} \rangle + \langle I_{\downarrow} \rangle$. On the other hand, we can consider the situation where $\Delta_{\mu_{\uparrow}} = -\Delta_{\mu_{\downarrow}}$, and thus there is no charge current through the conductor, $I_c = 0$. Still, there is a non-vanishing spin current $I_s \equiv \langle I_{\uparrow} \rangle - \langle I_{\downarrow} \rangle$, and according to the above result one can observe the current noise $S = e(1 - T) |I_s|$ induced by spin transport only.

We finally mention a further scenario where entanglement can be measured in the current [4]. A double-dot system which is weakly coupled to an ingoing (1) and an outgoing (2) lead held at chemical potentials $\mu_{1(2)}$, but where now an electron coming from lead 1 (2) has the option to tunnel into both dots 1 and 2 with amplitude $\Gamma$. This results in a closed loop, and applying a magnetic field, an Aharonov-Bohm phase $\varphi$ will be picked up by an electron traversing the double-dot. In the Coulomb blockade regime we find that due to cotunneling the current traversing the double-dot becomes (for $U > |\mu_1 \pm \mu_2| > J > k_B T, 2\pi \nu \Gamma^2$, with $U$ the single-dot charging energy)

$$I = e\pi \nu^2 \Gamma^4 \frac{\mu_1 - \mu_2}{\mu_1 \mu_2} (2 \pm \cos \varphi),$$

where the upper (lower) sign refers to triplet (singlet) state in the double-dot. Similarly, we find for the noise, $S_{\omega} \propto (2 \pm \cos \varphi)$. Thus, the current and noise reveal whether the double-dot is in a singlet or triplet state. The triplets can be further distinguished by applying spatially inhomogeneous magnetic fields leading to a beating of the AB phase oscillation due to the Berry phase [4].

References

1. A. Steane, Rept. Prog. Phys. 61, 117–173 (1998).
2. P. W. Shor, in Proc. 35th Symposium on the Foundations of Computer Science, (IEEE Computer Society Press), 124 (1994).
3. L. K. Grover, Phys. Rev. Lett. 79, 325 (1997).
4. D.P. DiVincenzo, D. Loss, cond-mat/9901137.
5. A. Barenco et al., Phys. Rev. A 52, 3457 (1995).
6. J. I. Cirac, P. Zoller, Phys. Rev. Lett. 74, 4091 (1995).
7. D. Cory, A. Fahmy, T. Havel, Proc. Nat. Acad. Sci. U.S.A. 94, 1634 (1997).
8. N. A. Gershenfeld, I. L. Chuang, Science 275, 350 (1997).
9. D. Loss, D. P. DiVincenzo, Phys. Rev. A 57, 120 (1998).
10. A. Shnirman, G. Schönh, Z. Hermon, Phys. Rev. Lett. 79, 2371 (1997).
11. D.V. Averin, Solid State Commun. 105, 659 (1998).
12. L. B. Ioffe et al., Nature 398, 679 (1999).
13. B. Kane, Nature 393, 133 (1998).
14. J. M. Kikkawa et al., Science 277, 1284 (1997); J.A. Gupta et al., Phys. Rev. B 59, R10421 (1999).
15. G. Burkard, D. Loss, D. P. DiVincenzo, Phys. Rev. B 59, 2070 (1999).
16. G. Burkard, D. Loss, E. Sukhorukov, cond-mat/9906071.
17. L. Jacak, P. Hawrylak, and A. Wójs, Quantum Dots (Springer, Berlin, 1997).
18. S. Tarucha et al., Phys. Rev. Lett. 77, 3613 (1996).
19. L. P. Kouwenhoven et al., Proceedings of the Advanced Study Institute on Mesoscopic Electron Transport, edited by L. L. Sohn, L. P. Kouwenhoven, G. Schönh (Kluwer, 1997).
20. R. C. Ashoori, Nature 379, 413 (1996).
21. R. J. Luyken, et al., preprint.
22. F. R. Waugh et al., Phys. Rev. Lett. 75, 705 (1995); C. Livermore et al., Science 274, 1332 (1996).
23. T. H. Oosterkamp et al., Phys. Rev. Lett. 80, 4951 (1998).
24. R. H. Blick et al., Phys. Rev. Lett. 80, 4032 (1998); ibid 81, 689 (1998). T. H. Oosterkamp et al., cond-mat/9809143.
25. C. H. Bennett and G. Brassard, in Proceedings of the IEEE International Conference on Computers, Systems and Signal Processing, Bangalore, India (IEEE, NY, 1984), p. 175.
26. A. Einstein, B. Podolsky, N. Rosen, Phys. Rev. 47, 777 (1935).
27. A. Aspect, J. Dalibard, G. Roger, Phys. Rev. Lett. 49, 1804 (1982); W. Tittel et al., Phys. Rev. Lett. 81, 3563 (1998).
28. D. Boumeester et al., Nature 390, 575 (1997); D. Boschi et al., Phys. Rev. Lett. 80, 1121 (1998).
29. G. Seelig, G. Burkard, D. Loss, preprint.
30. R. Loudon, in Coherence and Quantum Optics VI, eds. J. H. Eberly et al. (Plenum, New York, 1990).
31. A. Imamoglu et al., quant-ph/9904096.
32. R. Hanbury Brown and R. Q. Twiss, Nature (London) 177, 27 (1956).
33. M. Büttiker, Phys. Rev. Lett. 65, 2901 (1990); Phys. Rev. B 45, 12485 (1992).
34. T. Martin, R. Landauer, Phys. Rev. B 45, 1742 (1992).
35. R. C. Liu et al., Nature 391, 263 (1998); M. Henny et al., Science 284, 296 (1999); W. D. Oliver et al., Science 284, 299 (1999).
36. For a positive sign in the noise cross correlations due to the boson-like properties of Cooper pairs see, J. Torrèse, T. Martin, cond-mat/9906012.
37. G. A. Prinz, Science 282 1660 (1998).
38. For finite frequencies, we obtain the noise power
\[ S_{\alpha\alpha}(\omega) = (2e^2/h\nu)((1 - \delta_{\omega,0}) + 2T(1 - T)(\delta_{\omega,0} + \delta_{\omega,(1-e)(\nu)}))/h) \]
39. For instance, in metals such as bulk Cu the quasiparticle weight becomes within RPA approximation \( z_F \approx 0.77 \), while for a GaAs 2DEG we find \( z_F \approx 1 - r_s(1/2 + 1/\pi) = 0.66 \) for the GaAs interaction parameter \( r_s = 0.61 \).
40. V. A. Khilus, Zh. Eksp. Teor. Fiz. 93 2179 (1987).
41. D. Loss, E. Sukhorukov, preprint.