Entanglement of solid-state qubits by measurement

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We show that two identical solid-state qubits can be made fully entangled (starting from completely mixed state) with probability 1/4 just measuring them by a detector, equally coupled to the qubits. This happens in the case of repeated strong (projective) measurements as well as in a more realistic case of weak continuous measurement. In the latter case the entangled state can be identified by a flat spectrum of the detector shot noise, while the non-entangled state (probability 3/4) leads to a spectral peak at the Rabi frequency with the maximum peak-to-pedestal ratio of 32/3.

Prospective solid-state realizations of quantum computers have many potential advantages over other proposed realizations, including more natural scalability, simple electrical control of parameters, and use of very well developed technology. A number of theoretical proposals on solid-state quantum computers have been put forward (see, e.g. Ref. 8) and spectacular experimental results have been achieved, including demonstrations of charge qubits using single-Cooper-pair boxes, flux qubits using superconducting loops interrupted by Josephson junctions, and combined charge-flux qubits (the demonstrated quality factor was as high as 25,000 in Ref. 9). Obviously, next important experimental step is the demonstration of entangled solid-state qubits.

The qubits can be made entangled by their direct interaction if the interaction depends on the states of the qubits and at least one qubit is in a coherent superposition before the interaction. In this letter we will discuss an alternative way, when two solid-state qubits are made entangled just by their simultaneous measurement by one detector, which thus provides an indirect coupling between qubits. The procedure works with a probability less than unity, and in this respect it is somewhat similar to the operation of conditional quantum gates based on simple linear optical elements. Our procedure can also be thought of as an entanglement purification, in which instead of trading the number of qubit pairs for better entanglement, we trade the success probability while no initial entanglement is needed. One more analogy is the preparation of entangled atoms in an optical cavity by monitoring the cavity decay.

In contrast to qubits represented by photons, which are physically destroyed by the acts of measurement, solid-state qubits only change their state due to measurement, that allows somewhat more freedom in designing quantum operations. On the other hand, it is quite difficult to realize simple projective measurements of solid-state qubits because of typically weak coupling with detector. Therefore, instead of simple abrupt collapse, we have to deal with dephasing like processes in the case of ensemble measurements or with the continuous (weak) measurements in the case of single qubits.

The theory of non-averaged continuous measurement of single solid-state qubits is under active development during last 4 years and has already produced a number of experimental predictions (for some of them, see Ref. 10). One of the predictions is related to a direct continuous measurement of a qubit Rabi oscillations by a weakly coupled detector. The theoretical result is that the qubit oscillations can be evidenced by the peak of the detector current spectral density at the Rabi frequency; however, the peak height cannot be larger than 4 times the noise pedestal (this fact seems to have recent experimental confirmation 11).

In this letter we consider the case when two identical qubits perform Rabi oscillations, which are continuously measured by a detector equally coupled to two qubits. We will show that the system is gradually collapsed into one of the two regimes: either qubits become fully entangled (Bell state) that can be identified by absence of the spectral peak of the detector current, or qubits state fall into the orthogonal subspace that can be identified by the Rabi spectral peak, which for an ideal detector is 32/3 times higher than the noise pedestal. The probabilities of two scenarios are 1/4 and 3/4, respectively (starting from completely mixed state), so on average the peak-to-pedestal ratio is equal to 8, twice as large as for a single qubit.

Figure 1 shows possible realizations of our setup. In the first realization [Fig. 1(a)] each qubit is made of a double quantum dot (DQD), occupied by a single electron, while the detector is a quantum point contact (QPC) located in between DQDs. The second possible realization [Fig. 1(b)] is based on single-Cooper-pair-boxes (SCPB) as qubits, which are measured by a single-electron transistor (SET). Other possible realizations (not shown) can be based on flux qubits or combined charge-flux qubits.

In the Hamiltonian of the system,

\[ \mathcal{H} = \mathcal{H}_{QB} + \mathcal{H}_{DET} + \mathcal{H}_{INT}, \]  

(1)

the first term describes two qubits alone, \( \mathcal{H}_{QB} = (\varepsilon_a/2)(a_\uparrow^\dagger a_\downarrow - a_\downarrow^\dagger a_\uparrow) + H_a(a_\uparrow^\dagger a_\downarrow + a_\downarrow^\dagger a_\uparrow) + (\varepsilon_b/2)(b_\uparrow^\dagger b_\downarrow - b_\downarrow^\dagger b_\uparrow) + H_b(b_\uparrow^\dagger b_\downarrow + b_\downarrow^\dagger b_\uparrow), \)

where \( \varepsilon_a \) and \( \varepsilon_b \) are energy asymmetries, which are neglected in this letter, \( \varepsilon_a = \varepsilon_b = 0 \), the amplitudes \( H_a \) and \( H_b \) describe the tunneling within qubits (we assume \( H_a = H_b \) in the most of this letter), and the direct interaction term \( U a_\uparrow^\dagger a_\downarrow b_\downarrow^\dagger b_\uparrow \) is neglected. The frequencies of free Rabi oscillations of qubits, \( \Omega_a = \).
tor voltage is applied during short time intervals. Since and 
we use $\hbar = 1$) obviously coincide, $\Omega_a = \Omega_b = \Omega$ 
$[1, 2]$. For simplicity we limit ourselves by the case of DQD qubits, measured 
by a low transparency QPC (though generalization to other cases is simple), so that the detector Hamiltonian is 
$H_{DET} = \sum_{i} E_i c_i^\dagger c_i + \sum_{i, \sigma} V_i c_i^\dagger c_i + \sum_{i, \sigma} T_i (c_i^\dagger c_{\sigma} + c_{\sigma}^\dagger c_i)$ 
and the interaction term is $H_{INT} = \sum_{i, \sigma} \Delta T_a (a_i^\dagger a_{\sigma} - a_i^\dagger a_{\sigma}^\dagger) (c_{\sigma}^\dagger c_1 + c_1^\dagger c_{\sigma}) + \sum_{i, \sigma} \Delta T_b (b_i^\dagger b_{\sigma} - b_i^\dagger b_{\sigma}^\dagger) (c_{\sigma}^\dagger c_2 + c_2^\dagger c_{\sigma})$; we will be mostly interested in the case of equal coupling, 
$\Delta T_a = \Delta T_b$.

The four basis states of two qubits, $|1\rangle \equiv |a_b\rangle$, $|2\rangle \equiv |a_b\rangle$, $|3\rangle \equiv |a_b\rangle$, $|4\rangle \equiv |a_b\rangle$, correspond to 
4 values of the average current through the detector: 
$I_{1,2,3,4} = 2 \pi (T \pm \Delta T_a \pm \Delta T_b) \rho_{I,\sigma} e^2 V$, where $V$ is the 
QPC voltage and $\rho_{I,\sigma}$ are densities of states. It is important 
that in the case of equal coupling two currents coincide, 
$I_2 = I_3 \equiv I_{23}$, so the measurement cannot distinguish 
between states $|2\rangle$ and $|3\rangle$. Besides the “measurement” basis, it is convenient to introduce also the Bell basis: 
$|1\rangle_B \equiv (|a_b\rangle - |a_b\rangle)/\sqrt{2}$, $|2\rangle_B \equiv (|a_b\rangle - |a_b\rangle)/\sqrt{2}$, $|3\rangle_B \equiv (|a_b\rangle + |a_b\rangle)/\sqrt{2}$, and $|4\rangle_B \equiv (|a_b\rangle + |a_b\rangle)/\sqrt{2}$. Notice that $|1\rangle_B$ and $|2\rangle_B$ are eigenstates of 
$H_{QB}$ if $H_a = H_b$, while the states $|3\rangle_B$ and $|4\rangle_B$ are transformed by $H_{QB}$ as 
$\cos(2\Omega t + \phi)|3\rangle_B - \sin(2\Omega t + \phi)|4\rangle_B$.

Before considering continuous measurements, let us discuss a simpler case of a sequence of “orthodox” projective measurements which can be realized if the coupling with the detector is strong ($C \gg 1$, see below) and detector 
voltagoe is applied during short time intervals. Since the 
states $|2\rangle$ and $|3\rangle$ are mutually indistinguishable, the 
the two-qubit density matrix $\rho$ is projected each time into 
one of three subspaces, corresponding to states $|1\rangle$, $|23\rangle$, and $|4\rangle$ (we use notation $[23]$ for the subspace spanned by $|2\rangle$ and $|3\rangle$). The sequence of measurements separated by time periods $\Delta t$ can be described by such projections separated by intervals of unitary evolution due to $H_{QB}$.

Suppose the first measurement resulted in the current $I_{23}$, then the system is projected into $[23]$, which is also a subspace $[13]$ in the Bell basis. If the state would be exactly $|1\rangle_B$ (which does not evolve under $H_{QB}$), then all subsequent measurements would give the same result $I_{23}$ and the state $|1\rangle_B$ would remain unchanged. However, if the two qubits would be in the state $|3\rangle_B$, then the next measurement would result again in $I_{23}$ only with probability $p = (\cos \Omega \Delta t)^2$, while the probabilities of results $I_1$ and $I_4$ would be $(1-p)/2$ each. Therefore, if a sufficiently long sequence of current measurements repeatedly gives the result $I_{23}$, the two-qubit density matrix $\rho$ purifies and becomes close to the fully entangled state $|1\rangle_B$.

Simple analysis shows that after $N$ successful measurements (all results are $I_{23}$)

$$\rho_{11}(N) = \frac{\rho_{11}(0)}{\rho_{11}(0) + \rho_{33}(0) (\cos \Omega \Delta t)^{2(N-1)}}$$

where $\rho_{11}(0)$, $\rho_{33}(0)$, and $\rho_{11}(N)$ are the corresponding 
density matrix elements in the Bell basis before and after the sequence of measurements, while the probability 
of a successful sequence is $P(N) = \rho_{11}(0) + \rho_{33}(0) (\cos \Omega \Delta t)^{2(N-1)}$. For large $N$ the difference between 
the obtained state and state $|1\rangle_B$ becomes exponentially small, while the probability of success is close to $\rho_{11}(0)$, which is equal to $1/4$ for the fully mixed initial state $\rho_{11,\text{mix}} = \rho_{33,\text{mix}} = \delta_{ij}/4$ (this state is a direct product of fully mixed states of each qubit). The purification rate depends on $\Delta t$, and is the fastest when $\Delta t$ is close 
to $(1/4 + k^2)/2 \pi / \Omega$ ($k$ is an integer), which is a regime 
opposite to the quantum nondemolition measurement.

Notice that if some measurement in a sequence resulted 
in the current $I_1$ or $I_4$, then the subsequent measurements 
(for general $\Delta t$) can result in any current $I_1$, $I_4$, 
or $I_{23}$; however, $\rho_{11}$ remains exactly zero (therefore, long 
sequences of $I_{23}$ results become extremely improbable). 
Hence, to obtain the Bell state $|1\rangle_B$, one have to apply 
some perturbation which mixes two subspaces (for example, 
apply a noise which affects $e_a$ and/or $e_b$) and repeat 
the procedure. Assuming a well-mixed case, the probability 
$1 - (3/4)^M$ of obtaining the state $|1\rangle_B$ can be made arbitrary 
close to unity by allowing sufficiently large number $M$ of attempts.

The procedure can obviously be used for the preparation 
of entangled states in a solid-state quantum computer, so it is important to discuss what happens if the 
conditions $H_a = H_b$ and $I_2 = I_3$ are not satisfied exactly. In the case of slightly different $H_a$ and $H_b$, Eq. 
changes insignificantly (cos $\Omega \Delta t$ should be replaced 
with $\cos \Omega \Delta t / \cos 2^{-1} \Delta \Omega \Delta t$, where $\Delta \Omega \equiv \Omega_a - \Omega_b$), however, the probability of an $N$-long successful sequence becomes $P(N) = \rho_{11}(0) (\cos 2^{-1} \Delta \Omega \Delta t)^{2(N-1)} + \rho_{33}(0) (\cos \Omega \Delta t)^{2(N-1)}$ and 
decreases to zero at $N \to \infty$. Estimating the average length of a successful sequence, $\bar{N} \approx (\sin 2^{-1} \Delta \Omega \Delta t)^{-2}$, one can estimate a typical deviation from state $|1\rangle_B$,

$$1 - \rho_{11} B \sim (\cos \Omega \Delta t / \cos 2^{-1} \Delta \Omega \Delta t)^{2(\sin 2^{-1} \Delta \Omega \Delta t)^2},$$

which is as small as $\sim \exp[-(2\Omega / \Delta \Omega)^2]$ if $\Delta t \ll \Omega^{-1}$ (Quantum Zeno regime) and even smaller, $\sim (\cos \Omega \Delta t/(2\pi / \Omega \Delta t)^2)$, if $\Delta t$ is close to $\pi / 2\Omega$. 

FIG. 1. Schematic of two qubits measured by an equally 
coupled detector. (a): Realization based on double quantum 
dots measured by a quantum point contact, (b): realization 
based on single-Cooper-pair-boxes measured by a single- 
electron transistor. Measurement can entangle qubits.
To analyze the effect of a small difference between $I_2$ and $I_3$ because of slightly different coupling, we have to use the standard theory\cite{2,3} of weak quantum measurements and take into account the detector shot noise $S_t = 2eI_t$. We assume that during short measurement interval $\delta t$ the currents $I_1$ and $I_2$ can be unambiguously identified, while current $I_2$ and $I_3$ are almost indistinguishable and small corresponding signal-to-noise ratio is characterized by the parameter $\epsilon = (I_2 - I_3^2)/4D \ll 1$ where $D = S_{23}/2\delta t$ is the variance of the measured noisy current. Each successful measurement tends to shift the state towards either $|2\rangle$ or $|3\rangle$ and so decreases the amount of entangled state $|1\rangle_B$, that competes with the purification due to Eq. (2) and leads to iterative formula $\rho_{11}^B(N+1) \approx \rho_{11}^B(N) - \epsilon/4 + |1 - \rho_{11}^B(N)|^2/\sin\Omega\delta t)^2$ valid when $\rho_{11}^B$ is close to unity (exact formula is longer). Therefore, a typical deviation from pure entanglement $1 - \rho_{11}^B \approx \epsilon/4(\sin\Omega\delta t)^2$ scales as $\epsilon$.

Now instead of instantaneous measurements let us consider a more realistic case of a continuous measurement, realized when the detector voltage is applied all the time. For the analysis we will use the Bayesian formalism\cite{3} assuming weakly responding linear detecting regime, $|\Delta I_{a,b}| \ll I_i$, $\Delta I_a = I_1 - I_3 = I_2 - I_4$, $\Delta I_b = I_1 - I_2 = I_3 - I_4$, and concentrating on the case of symmetric weak coupling, $C_a \equiv C_b \ll 1$, $C_a = (\Delta I_a)^2/2S_0H_a$, $C_b = (\Delta I_b)^2/2S_0H_b$, where the frequency-independent detector shot noise spectral density $S_0$ does not depend significantly on the qubits state.

The evolution of the two-qubit density matrix $\rho$ can be described by the equation\cite{2,3} (in Itô representation)

$$\frac{d}{dt}\rho_{ij} = (I(t) - \sum_k \rho_{kk}I_k)(I_i + 1/2 - \sum_k \rho_{kk}I_k) \frac{\rho_{ij}}{S_0} - (I_i - I_j)^2/4S_0 + \gamma_{ij} \rho_{ij} - i[\mathcal{H}_{QB}, \rho]_{ij},$$

where the extra dephasing rate $\gamma_{ij} = (\eta^{-1} - 1)(I_i - I_j)^2/4S_0$ depends on detector identity $\eta (0 \leq \eta \leq 1)$ and vanishes for the QPC as a detector\cite{2,3} ($\eta = 1$); however, this term is important, for example, for the SET. To simulate individual realizations of the random measurement process, the noisy detector current $I(t)$ can be calculated as

$$I(t) = \xi(t) + \sum_k \rho_{kk}I_k,$$

where $\xi(t)$ is a white noise with spectral density $S_0$. Notice that the averaged dynamics (master equation) can be obtained by averaging over noise $\xi(t)$ that eliminates the first term in Eq. (4).

We performed extensive Monte Carlo simulations\cite{4} and found the following (Fig. 2). In the symmetric case, $H_a = H_b, C_a = C_b = C$ [$C \equiv (C_a + C_b)/2$; we mostly used $C$ between 1/4 and 1], any initial state either evolves eventually into the fully entangled Bell state $|1\rangle_B$ ($\rho_{11}^B \rightarrow 1$) or ends up in the orthogonal subspace ($\rho_{11}^B \rightarrow 0$) performing oscillations within this subspace so that the "signal" $z \equiv \rho_{11} - \rho_{44} = 2\Re\rho_{23}^B$ (which affects the detector current) oscillates with frequency $\Omega$ and fluctuating amplitude within the range from 0 to 1. Within the accuracy consistent with the number of trials, the probability of evolving into state $|1\rangle_B$ coincides with $\rho_{11}^B(0)$, similar to the case of orthodox measurement sequence. For $\eta = 1$ the two-qubit state eventually becomes pure independently of an initial choice (similar to the one-qubit case\cite{2,3}). This fact is obvious for the state $|1\rangle_B$, while in the oscillating scenario the surviving nondiagonal matrix elements in the Bell basis satisfy equations $(\Re\rho_{23}^B)^2 = \rho_{22}^B\rho_{33}^B$, $(\Im\rho_{23}^B)^2 = \rho_{22}^B\rho_{33}^B$, and $(\Im\rho_{23}^B)^2 = \rho_{22}^B\rho_{33}^B$. In the state $|1\rangle_B$ the numerically calculated spectral density of the detector current is flat and equal to $S_0$ (the signal $z$ is zero), while in the oscillating state it exhibits a peak (lower inset in Fig. 2) at frequency $\Omega$ with the peak height close to the analytical result $(32/3)S_0$ (see below).

The fact of collapsing eventually either into the state $|1\rangle_B$ or into the orthogonal subspace can be understood using an analogy with the sequential measurement case considered above, and is caused by the fact that neither unitary evolution due to $\mathcal{H}_{QB}$ nor nonunitary evolution due to measurement mixes two subspaces [see Eq. (3)]. The probability of two scenarios should be equal to the contribution of two subspaces to the initial density matrix, $\rho_{11}^B(0)$ and $1 - \rho_{11}^B(0)$, since ensemble averaged value $\langle \rho_{11}^B(t) \rangle$ does not change with time (as follows from the master equation).

To find analytically the spectral density of the detector current for the oscillating state, we have used two approaches\cite{4} leading to the same result. The first one is based on the master equation and collapse ansatz. Using classical equation $I(t) = z\Delta I + \xi(t)$, we calculate the current correlation function $K_f(\tau) = \langle I(0)I(\tau) \rangle$ as $K_f(\tau > 0) = (\Delta I)^2K_{\xi}(\tau)$, while $K_{\xi}(\tau)$ is calculated in the following way. At time $\tau = 0$ the two-qubit state is collapsed into one of the three basis states of the sub-

\begin{figure}[h]
\centering
\includegraphics[width=0.8\textwidth]{fig2}
\caption{Two numerical realizations of $\rho_{11}^B$ evolution starting from the fully mixed state. The upper (solid) line illustrates the scenario of collapse into fully entangled Bell state $|1\rangle_B$, while the lower (dotted) line shows a collapse into the orthogonal subspace. Two insets show the corresponding spectral densities $S_f(\omega)$ of the detector noise (solid/dashed lines are the numerical/analytical results).}
\end{figure}
with probability weights 1/3 each, we obtain $K \equiv \pm$ in the third one. Summing 3 contributions to $\langle z(0)z(\tau) \rangle$ with probability weights 1/3 each, we obtain $K_z = (2/3)G(\tau)$ and the current spectral density

$$S_I(\omega) = S_0 + \frac{8}{3} \frac{\Omega^2(\Delta I)^2 \Gamma}{(\omega^2 - \Omega^2)^2 + \gamma^2 \omega^2}. \quad (7)$$

In the case $\Gamma \ll \Omega$ the spectral peak at the Rabi frequency $\Omega$ corresponds to the $Q$-factor of $8\eta/C$ (similar to one-qubit case) and has the peak height equal to $(32/3)\eta S_0$, confirming the numerical result (see lower inset in Fig. 3).

The second method of $S_I(\omega)$ calculation is based on the Bayesian Eq. (2) assuming $\eta = 1$ and random evolution of a pure state with $z = A(t) \cos(\Omega t + \Phi(t))$ [then $y = A(t) \sin(\Omega t + \Phi(t))$]. In this method the correlation between noise $\xi(0)$ and evolution of the density matrix at later time should be taken into account, so $K_I(\tau > 0) = (\Delta I)^2 K_z(\tau) + \Delta K_{\xi z}(\tau)$, while correlation functions $K_z(\tau)$ and $K_{\xi z}(\tau)$ should be calculated by averaging of a long individual realization over time. We have proved that the result for $K_I(\tau)$ calculated by this method coincides with the result of the previous method for arbitrary coupling $C$; however, the formalism is much simpler for weak coupling, $C \ll 1$. In this case the stochastic differential equations for $A(t)$ and $\Phi(t)$ can be averaged over oscillations with frequency $\Omega$ and the correlation functions can be calculated analytically:

$$K_z(\tau) = (5/12)G(\tau) \text{ and } K_{\xi z}(\tau > 0) = G(\tau) \Delta I/4.$$  

This gives us a natural partition of the relative spectral peak height on 3/2 into two contributions: “classical” part 20/3 comes from oscillations of the signal $z$, while the “quantum” contribution equal to 4 is due to partial collapse of $\rho$ correlated with the detector noise. Comparing this partition with the partition $4 = 2 + 2$ for a one-qubit measurement, we see that the classical part grows faster than the quantum part when the number of qubits is increased.

Numerical simulations show that if the two Rabi frequencies $\Omega_a$ and $\Omega_b$ are slightly different, or small difference between $\xi_a$ and $\xi_b$ is due to asymmetry of coupling (different $\Delta I_a$ and $\Delta I_b$), then the two-qubit density matrix $\rho$ makes rare abrupt jumps between a state very close to $|1\rangle^B$ and the oscillating state. To find the switching rate analytically, we have used the master equation starting from entangled initial condition $\rho_{11} = 1$ and calculated the linear term in $\rho_{11}(t)$ dependence at $t \gg \Gamma^{-1}$ (but when $\rho_{11}(t)$ is still close to unity). In this way we have obtained the rate $\Gamma_{B\rightarrow O} = (\Delta \Omega)^2/2\Gamma$ of switching from the Bell state to the oscillating state due to slightly different Rabi frequencies, and the rate $\Gamma_{O\rightarrow B} = (\Delta \epsilon/C)^2/8\gamma$ when $\Omega_a = \Omega_b$, but couplings $\Delta_a$ and $\Delta_b$ are slightly different. To find the rate of the reverse switching, notice that the stationary master equation has the solution $\rho_{11, st} = 1$ and calculated the linear term in $\rho_{11, st}(t)$ dependence at $t \gg \Gamma^{-1}$, and so $\Gamma_{O\rightarrow B} = \Gamma_{B\rightarrow O}/3$. The numerical histograms of switching time distributions confirm these formulas. Taking into account rare switching events, the average spectral density of the detector current is given by Eq. (6) multiplied by 3/4, so the spectral peak height is equal to $8\eta S_0$.

Finally, we have studied the effect of environmental dephasing, modeling it with two small dephasing rates $\gamma_a$ and $\gamma_b$ acting separately onto two qubits. This leads to a slightly mixed state, even for an ideal detector and to switching events (similar to the case above) with $\Gamma_{B\rightarrow O} = 3\gamma_a/\gamma = (\gamma_a + \gamma_b)/2$. Notice that a controllable weak external noise can be used in a simple feedback protocol to restore the entangled state after undesirable switching to the oscillating state.

In conclusion, we have found that the continuous measurement of two identical solid-state qubits by equally coupled detector leads to either full spontaneous entanglement of qubits (Bell state) or to collapse into orthogonal oscillating state. In the latter case the noisy detector current has a spectral peak at the Rabi frequency, while in the former case the spectrum is flat. Slight asymmetry of the two qubit configuration as well as environmental dephasing leads to switching events leading to two regimes. It is important to mention that for an experimental observation of the phenomenon the quantum ideality $\eta$ of the detector should not necessarily be close to unity; it should only be large enough to allow distinguishing the Rabi spectral peak with the peak-to-pedestal ratio of $32\eta/3$.

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28. Notice that if $z$ would oscillate with the maximum possible amplitude of 1, the corresponding contribution to the peak height would be equal to $8S_0$, which is still less than $32S_0/3$. 