Operation of a phase qubit as a quantum switch

Jian Li,[1,*] G. S. Paraoanu,[1] Katarina Cicak,[2] Fabio Altomare,[2,†] Jae I. Park,[2] Raymond W. Simmonds,[2] Mika A. Sillanpää,[1] and Pertti J. Hakonen[1]

[1] *Low Temperature Laboratory, Aalto University, PO Box 15100, FI-00076 AALTO, Finland*
[2] *National Institute of Standards and Technology, 325 Broadway, Boulder, Colorado 80305, USA*

We present an experimental demonstration of a phase qubit acting as an on/off switch for the absorbtion of photons in a probe microwave beam. The switch is controlled by a second control microwave field. The on/off states of the qubit are steady states which for suitable values of the control field can be close enough to dark states. The functioning of the device is based on the Autler-Townes effect. A full numerical model up to the fifth energy level, including independently-determined parameters for relaxation, dephasing, and cross-couplings, gives excellent agreement with the experiment and allows us to characterize the time-scale for switching.

Routers, switches, and repeaters are some of the essential components of modern information-processing architectures. Similar devices will be needed in future quantum computers, possibly forming the backbone of a quantum internet [1]. In the optical regime, quantum switches have been demonstrated with a single ^{87}Rb atom in a high-finesse cavity [2], and also as heralded single-photon absorbtion by one trapped ion [3]. The functioning of our device makes use of the first three levels of a phase qubit and of the so-called Autler-Townes effect [4], which has been demonstrated recently for superconducting qubits [5–7]. With improved dissipation and decoherence, in such systems it is possible to demonstrate directly electromagnetic-induced transparency effects [8], which opens up to other applications using the physics of three levels - for example single-atom amplifiers [9].

In this paper we show how the phase qubit can be used as a device which absorbs or not an itinerant (flying) qubits - in our case realized by microwave photons - depending on weather a control (coupling) field at a different frequency is shone on the qubit. We demonstrate that the state of the qubit can be flipped between two stationary states which are defined by the strength of the control field. For suitable values of the control field, these states approximate ideal dark states.

Our device has been described elsewhere [6]. It consists of a phase qubit (a rf SQUID with loop inductance L, junction capacitance C and Josephson energy E_J). The device forms a multilevel system with the first three levels denoted by $|0\rangle$, $|1\rangle$, and $|2\rangle$. The Hamiltonian of the phase qubit is

$$H = \frac{Q^2}{2C} + \frac{(\Phi - \Phi_{\rm ext})^2}{2L} - E_J \cos\left(2\pi\frac{\Phi}{\Phi_0}\right), \qquad (1)$$

where Φ is the flux (phase) variable and $\Phi_{\rm ext} = \Phi_{\rm dc} + \Phi_{\rm rf}(t)$ is the total magnetic-flux component of the externally applied d.c. and r.f. fields. In this system, the r.f. radiation with frequencies close to the resonance of the two transitions $|0\rangle \to |1\rangle$ and $|1\rangle \to |2\rangle$ can be applied via coplanar-waveguide transmission lines. The Rabi frequencies are denoted by Ω_p and Ω_c and they are directly proportional to the amplitude of the r.f. fields. Dissipation processes are included in the model by the standard Markov-approximation Liouville superoperators, with relaxation part

$$\mathcal{L}_{rel}[\rho] = \frac{\Gamma_{10}}{2}\left(2\sigma_{01}\rho\sigma_{10} - \sigma_{11}\rho - \rho\sigma_{11}\right) + \frac{\Gamma_{21}}{2}\left(2\sigma_{12}\rho\sigma_{21} - \sigma_{22}\rho - \rho\sigma_{22}\right) + \kappa\left(\sigma_{01}\rho\sigma_{21} + \sigma_{12}\rho\sigma_{10}\right), \qquad (2)$$

and dephasing

$$\mathcal{L}_{dep}[\rho] = \frac{\Gamma_1^{\varphi}}{2}\left(2\sigma_{11}\rho\sigma_{11} - \sigma_{11}\rho - \rho\sigma_{11}\right) + \frac{\Gamma_2^{\varphi}}{2}\left(2\sigma_{22}\rho\sigma_{22} - \sigma_{22}\rho - \rho\sigma_{22}\right). \qquad (3)$$

Here ρ is the density matrix in the Schrödinger picture, the interlevel relaxation rates between $|1\rangle \to |0\rangle$ and $|2\rangle \to |1\rangle$ are denoted as Γ_{10} and Γ_{21}, respectively, and Γ_1^{φ}, Γ_2^{φ} are the dephasing rates of the states $|1\rangle$ and $|2\rangle$. Also, $\kappa = \sqrt{\Gamma_{10}\Gamma_{21}}$, $\sigma_{ij} \stackrel{\rm def}{=} |i\rangle\langle j|$. The dissipation parameters were either measured directly or extracted from spectroscopy data [6, 10], and for this sample they were $\Gamma_{21} \approx 2\pi \times 11$ MHz, $\Gamma_{10} \approx 2\pi \times 7$ MHz, $\Gamma_1^{\varphi} \approx 2\pi \times 7$ MHz, $\Gamma_2^{\varphi} \approx 2\pi \times 16$ MHz.

With these notations, the Markovian master equation for the density matrix in the Schrödinger picture ρ takes the form

$$\dot{\rho} = -\frac{i}{\hbar}[H, \rho] + \mathcal{L}_{rel}[\rho] + \mathcal{L}_{dep}[\rho]. \qquad (4)$$

This equation can be studied both analytically and numerically in the stationary state [10], with excellent agreement with the the experimental data [6].

Here we demonstrate that three-level atoms can be operated as well in the time-domain. As such, the phase qubit can be used as a switchable absorbant, and the

*Electronic address: jianli@ltl.tkk.fi
†Present affiliation: D-Wave Systems Inc., 100-4401 Still Creek Drive, Burnaby, British Columbia V5C 6G9, Canada

switching times for our system are determined experimentally. In our experiment, the intensities of the probe field are such that only one photon is present in the system during the initialization-to-measurement period. Thus the device can be used to allow or block the transmission of single photons. In Fig. 1 we present the pulse sequence used to demonstrate the functioning of the three-level quantum switch time-domain based on the Autler-Townes effect. The Hamiltonian Eq. (1) is dynamically manipulated by using the coupling (pump) field, coupling resonantly the levels $|1\rangle$ and $|2\rangle$. The pump is pulsed starting at 150 ns and ending at 300 ns. To probe the system we use a continuous signal, with frequency around the transition $|0\rangle \rightarrow |1\rangle$. A standard measure pulse follows, allowing us to record the probability $P_{|1\rangle} + P_{|2\rangle}$ that the system is either one of the excited states $|1\rangle$ and $|2\rangle$.

In Fig. 2 we show the results of the actual experiment. In both the on-set phase as well as in the switch-off stage of the effect, the timescale for the system to reach the steady state is of the order of 50 ns, in qualitative agreement with the decoherence times in our system.

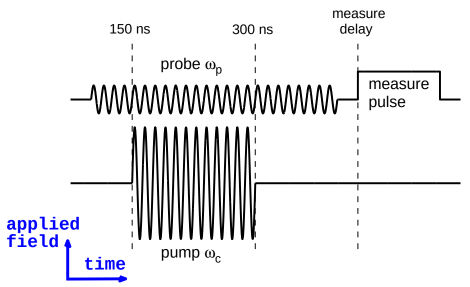

FIG. 1: Schematic of the pulse sequence for the dynamic Autler-Townes effect.

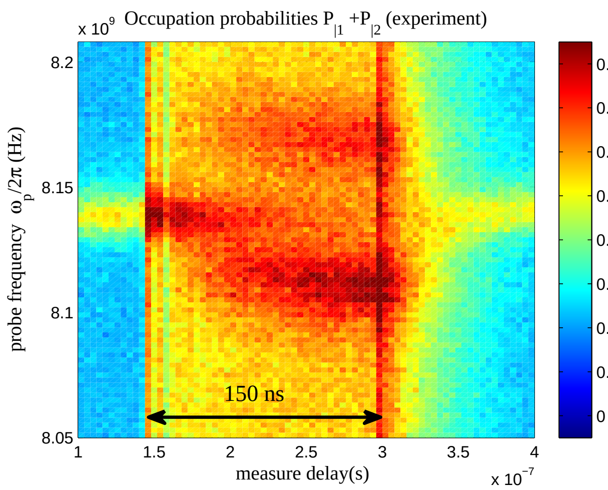

FIG. 2: On-set and extinction of the Autler-Townes effect. The figure shows the experimentally determined sum of excited-states level-occupancy probabilities, $P_{|1\rangle} + P_{|1\rangle}$.

In Fig. 3 we present the results of numerical simulations using Eq. (4). In the absence of the control field, the state $|2\rangle$ is empty, while the state $|1\rangle$ is populated when the probe field is resonant with the $|0\rangle \rightarrow |1\rangle$ transition. The agreement with the experiment is very good. For these simulations, the amplitudes of the probe and coupling fields correspond to Rabi frequencies $\Omega_p = 2\pi \times 3.45$ MHz (at the frequency of the first transition $\omega_{21} = 2\pi \times 7.975$ GHz), $\Omega_c = 2\pi \times 50$ MHz when the coupling field was on and a much smaller value $\Omega_c = 2\pi \times 1$ MHz when the coupling field was off. (The calibration of the field amplitudes is done in independent Rabi-frequency measurements between levels $|0\rangle$ and $|1\rangle$, and between levels $|1\rangle$ and $|2\rangle$.)

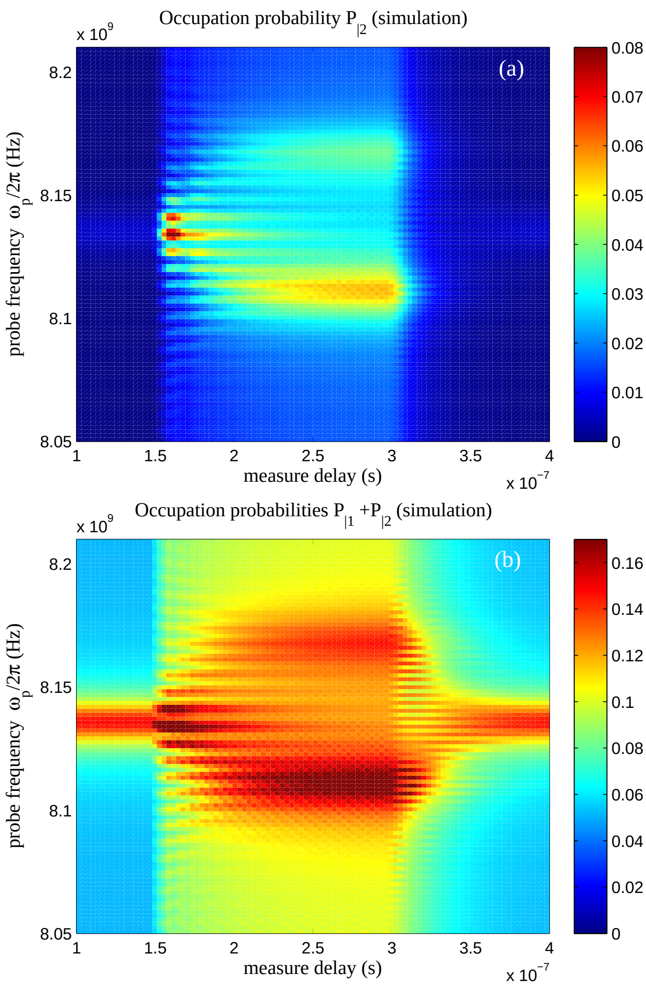

FIG. 3: Simulated occupation probabilities for the occupation probability of state $|2\rangle$ (a) and the sum of the occupation probabilities for the states $|1\rangle$ and $|2\rangle$ (b). The last one is the quantity measured directly in the experiment (see Fig. (2)).

The device can be used not only to absorb/transmit a photon but also as a switch of the quantum steady states of the qubit, which could have applications for quantum computing (see *e.g.* [11]). In our case, the switching occurs between two steady-states, $\rho_{\rm ON}^{(\rm st)}$ and

$\rho^{(st)}_{OFF}$. The ON state is characterized by the suppression of excitations to the state $|1\rangle$ normally caused by absorbtion of photons from the probe field. It is produced by a large value of the control field amplitude (effective Rabi frequency Ω_c). The state OFF corresponds to a lower value of the coupling field, which does not suppress excitations to the state $|1\rangle$. The steady-states are 3x3 density matrices which are found by solving numerically the equation $d/dt\rho^{(st)} = 0$. In the standard doubly-rotating frame corresponding to the two transitions [10] the Hamiltonian of this problem has eigenstates with zero eigenvalue (ideal dark states) with the structure $|D\rangle = (1/\sqrt{\Omega_c^2 + \Omega_p^2})(\Omega_c|0\rangle - \Omega_p|2\rangle$. For $\Omega_p \ll \Omega_c$ the dark state is approximately the ground state: if the system is trapped in such a state, the population of the excited levels is significantly suppressed. In the same limit, it can be shown that the steady state of the qubit becomes close to a dark state [10], with relative fidelity [12] $\mathcal{F}[|\mathrm{D}\rangle, \rho^{(st)}] = \sqrt{\langle \mathrm{D}|\rho^{(st)}|\mathrm{D}\rangle}$ approaching the unit.

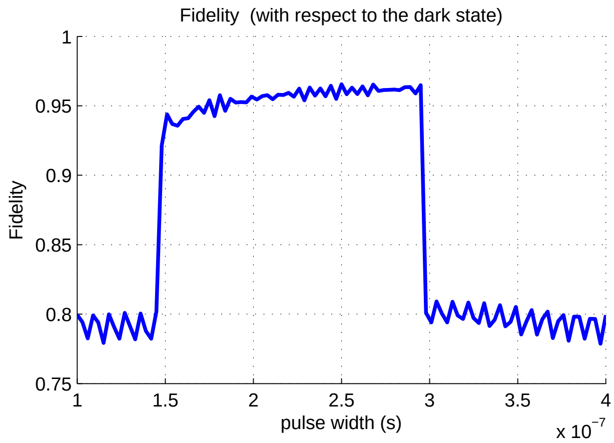

FIG. 4: Fidelity of the ON and OFF states with respect to the corresponding dark states. As expected, the ON state has a high fidelity, while the OFF state has a lower fidelity due to the use of a low-intensity coupling field.

In Fig. 4 we present numerical simulations for the fidelities in the ON and OFF state corresponding to the actual experiment. As mentioned above, the fidelity of the OFF steady-state ($\rho^{(st)}_{OFF}$) is expected to be lower than that of the ON state, reflecting the fact that the state $|1\rangle$ is populated (see Fig. 5).

In conclusion, we demonstrated the use of a phase qubit as quantum switch between two states with different properties; the device is operated in a fast, controllable way.

Acknowledgments

We acknowledge financial support from the Academy of Finland (nos. 129896, 118122, 130058, and 135135), from the National Graduate School of Material Physics, and from the European Research Council (StG).

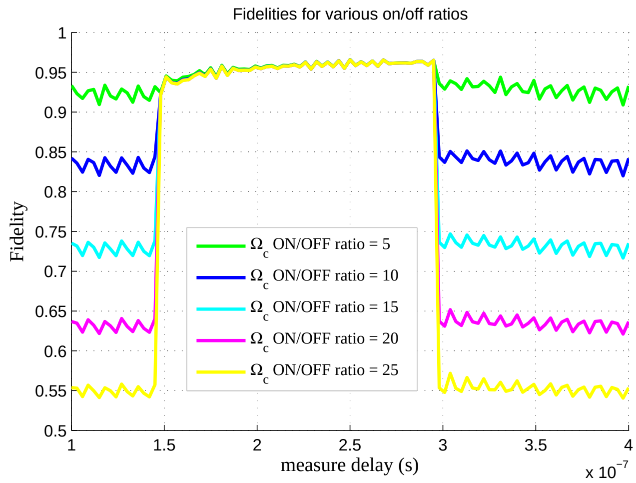

FIG. 5: Fidelities with respect to the respective dark states as a function of the on/off ratio of the coupling field.

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