Quantum dynamics of a dc-SQUID coupled to an asymmetric Cooper pair transistor

A. Fay\textsuperscript{1}, W. Guichard\textsuperscript{1}, O. Buisson\textsuperscript{1} and F. W. J. Hekking\textsuperscript{2}

\textsuperscript{1}Institut Néel, C.N.R.S.- Université Joseph Fourier, BP 166, 38042 Grenoble-cedex 9, France and
\textsuperscript{2}LPMMC, C.N.R.S.- Université Joseph Fourier, BP 166, 38042 Grenoble-cedex 9, France

We present a theoretical analysis of the quantum dynamics of a superconducting circuit based on a highly asymmetric Cooper pair transistor (ACPT) in parallel to a dc-SQUID. Starting from the full Hamiltonian we show that the circuit can be modeled as a charge qubit (ACPT) coupled to an anharmonic oscillator (dc-SQUID). Depending on the anharmonicity of the SQUID, the Hamiltonian can be reduced either to one that describes two coupled qubits or to the Jaynes-Cummings Hamiltonian. Here the dc-SQUID can be viewed as a tunable micron-size resonator. The coupling term, which is a combination of a capacitive and a Josephson coupling between the two qubits, can be tuned from the very strong- to the zero-coupling regimes. It describes very precisely the tunable coupling strength measured in this circuit\textsuperscript{1} and explains the ‘quantonium’ as well as the adiabatic quantum transfer read-out.

1. INTRODUCTION

Two quantum systems with discrete energy levels coupled to each other form an elementary block, useful for the study of fundamental phenomena and effects in quantum physics, especially in the context of quantum information. The interaction between the two quantum systems is essential to implement important concepts in this field such as entanglement, quantum gate operations, and quantum information transfer. As to the theoretical description of interacting quantum systems, two problems have been extensively studied in particular: a two-level system (or qubit) coupled to a harmonic oscillator and two coupled qubits. The former was used to describe among others the quantum electro-dynamics associated with atoms in a cavity\textsuperscript{2}, trapped ions coupled to their vibrations\textsuperscript{3} and more generally interaction between matter and photons\textsuperscript{4}. The latter was developed to describe entangled photons, trapped ions, and two-qubit quantum gate operations\textsuperscript{5}.

These studies were experimentally realized in the fields of quantum optics and atomic physics and more recently extended to include quantum solid state devices. In particular, during the last decade superconducting circuits have demonstrated their potential in connection with quantum experiments\textsuperscript{6,15}. They now appear as experimental model systems for studying fundamental quantum physics and basic blocks for quantum information.

In this paper we consider a superconducting circuit composed of two well-known elements coupled to each other: an inductive dc-SQUID and a Cooper pair transistor. This circuit constitutes an elementary building block that can be operated in various parameter regimes characterized by different types of quantum dynamics, as has been shown experimentally in the past. As we will detail below, this is possible due to the fact that three strongly coupled quantum variables determine the dynamics of this coupled circuit.

For instance, when the current-biased dc-SQUID is non-inductive and classical and the transistor is symmetric, one recovers the Quantronium\textsuperscript{10}. When the SQUID is operated in the nonlinear regime, the resulting system consists of a charge qubit coupled to an anharmonic oscillator; this system has been shown to allow for non-destructive quantum measurements\textsuperscript{16}. We recently operated the SQUID in the quantum few-level limit\textsuperscript{1}, demonstrating a very strong tunable coupling between two different types of qubit: a phase qubit and a charge qubit.

The experimental performance of this circuit was limited by uncontrolled decoherence sources. However its integration (three quantum variables strongly coupled on a micrometer scale), its tunability, its various optimal points make this circuit attractive once decoherence sources will be overcome upon technological improvements.

In this paper we present a rigorous theoretical analysis of the circuit in the parameter range of our experiments\textsuperscript{1}. The full Hamiltonian of the coupled circuit is presented, describing a two-level system (Cooper pair transistor) coupled to two anharmonic oscillators (dc-SQUID). In the relevant parameter range, the dc-SQUID dynamics can be reduced from two-dimensional to one-dimensional. Consequently the dynamics of the circuit is that of a qubit coupled to a single anharmonic oscillator. Depending on the anharmonicity, different regimes can be studied in this unique circuit. When the anharmonicity is neglected, we recover the physics of a qubit coupled to a harmonic oscillator. Its quantum dynamics can then be described by the well-known Jaynes-Cummings Hamiltonian. Although this limit can also be achieved with a qubit coupled to a high-Q coplanar wave guide cavity\textsuperscript{6,18} we wish to emphasize that the use of a dc-SQUID is of interest as it constitutes a micron-size resonator and it is tunable. When the oscillator is considered anharmonic, its interaction with a qubit gives rise to very complex dynamics which has been very little explored. If only the two lowest levels of the oscillator are considered, it can be reduced to a two-level system. The coupled circuit then describes two interacting qubits. In addition to the possibility to study different dynamical regimes depending on the anharmonicity of the resonator, the coupling between the SQUID and the transistor is
fully tunable. As a result the system can be operated at zero coupling, as well as in the weak and in the strong coupling limits.

In section II, after a description of the circuit under consideration, we construct the Lagrangian and the Hamiltonian using Devoret’s circuit theory [19]. Section III is devoted to the two-dimensional dynamics characteristic for an inductive SQUID and its reduction to one-dimensional dynamics in the relevant parameter range. In section IV we discuss the quantum dynamics of the asymmetric Cooper pair transistor, especially at its two optimal points where it is insensitive to noise fluctuations. In section V the terms describing the coupling between the dc-SQUID and the transistor are discussed. In section VI, the full quantum dynamics of the coupled circuit is presented. There we also discuss the two possible quantum measurements of the charge qubit that can be performed by the dc-SQUID. The last section discusses the tunable coupling strength of the circuit and its comparison with the experiments.

II. COUPLED CIRCUIT DYNAMICS

A. Circuit description

A schematic electronic representation of the circuit studied theoretically hereafter is presented in Fig. 1(b). In this circuit two different elements are coupled that correspond to two basic blocks for typical superconducting quantum devices. The first element is a dc-SQUID: it corresponds to a loop which contains two identical Josephson junctions (JJ), each with a critical current $I_0$ and a capacitance $C_0$. The total inductance $L_S$ of dc-SQUID is the sum of three inductances $L_1$, $L_2$ and $L_3$ associated with the different parts of the SQUID loop. The second element of the circuit is an asymmetric Cooper Pair transistor (ACPT) which is placed in parallel with the dc-SQUID. The ACPT consists of a superconducting island connected to the dc-SQUID by two different Josephson junctions. We denote by $I_1^T$ and $I_2^T$ the critical currents and $C_1^T$ and $C_2^T$ the capacitances of these junctions. The asymmetry of the transistor is characterized by the Josephson asymmetry parameter $\mu = (I_2^T - I_1^T)/(I_2^T + I_1^T)$ and the capacitance asymmetry parameter $\lambda = (C_2^T - C_1^T)/(C_2^T + C_1^T)$. The ACPT is also coupled to a gate-voltage; this is modeled theoretically by an infinite capacitance $C_P$ with a charge $Q_b$ with $C_P, Q_b \to \infty$ so the ratio $Q_b/C_0 \equiv V_g$. The circuit is current-biased, modeled similarly with the help of an infinite inductance $L_b \to \infty$ threaded by a flux $\Phi_b \to \infty$ so that the ratio $\Phi_b/L_b \equiv I_b$ remains constant. The properties of the circuit depend on four, experimentally tunable parameters $V_g$, $I_b$ and the fluxes $\Phi_S$ and $\Phi_T$ threading the dc-SQUID loop and the other loop of the circuit, respectively. As we will see, these parameters allow to control and change the dynamics of the coupled circuit. This circuit, realized and studied experimentally by A. Fay et al. [41], is shown in the SEM view in Fig. 1(a).

B. Devoret’s circuit theory

The relevant degrees of freedom of a superconducting circuit and their dynamics can be determined using the concept of node phases introduced by M. Devoret [19]. We distinguish two different kinds of nodes in the circuit. We first choose a ground node to which the zero of phase is associated. Note that this choice corresponds to a choice of gauge and is therefore arbitrary; although this choice affects the detailed form of the Hamiltonian it does not affect the resulting dynamics of the circuit. The other nodes of the circuit are called the active nodes, each described by an active phase. Six active nodes $\varphi_1, \varphi_2, \psi, \theta, \gamma$ and $\xi$.

FIG. 1. The coupled circuit studied experimentally in Ref. [1]. (a) On the right, a SEM picture shows the two elements of the circuit, i.e., an Asymmetric Cooper Pair Transistor (red frame) connected in parallel to a dc-SQUID (blue frame). A zoom of the transistor is shown on the left of the SEM picture. The transistor asymmetry stems from the difference between the areas of the two Josephson junctions. The quantum dynamics of the circuit can be modified by the bias current $I_b$, the flux $\Phi_g$ in the dc-SQUID, the flux $\Phi_T$ in the other loop and the gate-voltage $V_g$. (b) Schematic representation of the circuit. The four Josephson junctions are represented by a pure Josephson element in parallel with a capacitance. The voltage bias of the transistor gate and the current bias of the circuit are schematically indicated by an infinite capacitance $C_P$ and an infinite inductance $L_P$ respectively. The spanning tree drawn in red connects the ground node to the six active nodes $\varphi_1, \varphi_2, \psi, \theta, \gamma$ and $\xi$. Let us now introduce the so-called spanning tree [19]. Starting from the ground node we draw the branches of the spanning tree reaching each active node via a unique
path. In the case of the coupled circuit, the spanning tree (drawn in red in Fig. 1(b)), connects the ground node (bottom node) to the six active nodes.

The superconducting phase difference across a dipole in a superconducting circuit can be written as a function of the node phases with the help of two rules. We will illustrate these rules with the help of the circuit presented in Fig. 2 as an example. Here, three dipoles are placed in a loop threaded by the flux \( \Phi \). As a first rule, the superconducting phase difference across a dipole located on the spanning tree is given by the difference of the phases of the nodes linked to the dipole. Hence, the superconducting phases \( \Lambda_1 \) and \( \Lambda_3 \) of the dipoles 1 and 3, respectively, are given by \( \Lambda_1 = \phi_A \) and \( \Lambda_3 = \phi_B \). As a second rule, in the case of a dipole which is not located on the spanning tree, we first define a minimal loop which contains the previous dipole and the other dipole located on the spanning tree. The superconducting phase difference across the dipole is calculated, using the quantization of the phase for the minimal loop [20]. Let us apply this rule to determine the superconducting phase difference \( \Lambda_2 \) across dipole 2. The minimal loop corresponds simply to the loop of the circuit. The phase quantization in this loop gives \( \Lambda_2 = \phi_B - \phi_A - 2\pi \Phi/\Phi_0 \), with \( \Phi_0 \) the superconducting flux quantum. With the help of the two previous rules, we have analyzed the coupled circuit and expressed the superconducting phase difference across each dipole as a function of the six active phases.

\[ \Lambda_2 = \phi_B - \phi_A - 2\pi \Phi/\Phi_0 \]

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**FIG. 2.** Three dipoles in a loop threaded by the flux \( \Phi \). The spanning tree in red connects the ground node to the two active nodes \( \phi_A \) and \( \phi_B \). The superconducting phases \( \Lambda_1, \Lambda_2 \) and \( \Lambda_3 \) across the three dipoles respectively are functions of \( \phi_A, \phi_B \) and \( \Phi \).

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**C. Current conservation and system dynamics**

In a superconducting circuit, we find generally three different dipoles: a capacitance \( C \), an inductance \( L \) and a Josephson element \( J \). The current through each of these dipoles can be expressed as a function of the superconducting phase difference \( \Lambda \) across the dipole. The voltage \( V \) is related to the phase difference \( \Lambda \) by \( V = \phi_0 \Lambda \).

The current \( I_C \) through the capacitance is then given by \( I_C = \phi_0 C \Lambda \), where \( \phi_0 = \Phi_0/(2\pi) \). The current \( I_L \) through the inductance reads \( I_L = \phi_0 \Lambda / L \). Finally, the current through the Josephson element is given by the Josephson relation \( I_J = I_0 \sin(\Lambda) \), where \( I_0 \) is the Josephson critical current [21]. Using these expressions for the current, we can write the six current conservation laws for each active node in the studied circuit (the sum of the currents flowing into each node is zero). The six conservation laws yield six equations for the dynamics of the node phases (cf. appendix A), whose solution yields the dynamics of the entire circuit.

**D. The Lagrangian**

The Lagrangian \( \mathcal{L} \) of the circuit depends on the six node phases and their time derivatives. Let us define the vector \( \vec{x} = (\varphi_1, \varphi_2, \psi, \theta, \gamma, \xi) \) formed by the six node phases of the circuit. The six Euler-Lagrange equations are defined by [22]:

\[ \frac{d}{dt} \left( \frac{\partial \mathcal{L}}{\partial \dot{\vec{x}}} \right) - \frac{\partial \mathcal{L}}{\partial \vec{x}} = 0. \]  

The Lagrangian of the circuit has to be constructed in such a way that the Euler-Lagrange equations are equivalent to the current conservation equations (Appendix A). We take the Lagrangian to be of the following form:

\[ \mathcal{L}(\vec{x}, \dot{\vec{x}}) = T(\dot{\vec{x}}) - V(\vec{x}), \]  

with the kinetic energy

\[ T(\dot{\vec{x}}) = \phi_0^2 \left\{ \frac{1}{2} C_0 \dot{\varphi}_1^2 + \frac{1}{2} C_0 \dot{\varphi}_2^2 + \frac{1}{2} C_2^T (\dot{\xi} - \dot{\psi})^2 \right\} \]

\[ + \frac{1}{2} C_1^T \dot{\psi}^2 + \frac{1}{2} C_p (\dot{\theta} - \dot{\gamma})^2 + \frac{1}{2} C_p \dot{\theta}^2 \]  

and the potential energy

\[ V(\vec{x}) = -\phi_0 \left\{ I_0 \cos(\varphi_1) + I_0 \cos(\varphi_2) + I_1^T \cos(\xi - \psi) + I_1^T \cos(\psi) \right\} \]

\[ + \dot{\phi}_0^2 \left\{ \frac{1}{2} (\xi - \varphi_1)^2 \right\} L_1 + \frac{1}{2} \left( \gamma - \varphi_2 \right)^2 \]

\[ + \frac{1}{2} \frac{1}{L_3} \left( \gamma - \xi - \phi_S \right)^2 - \frac{1}{\phi_0} I_p \gamma \].

The kinetic energy corresponds to the energy stored in the capacitances of the circuit, whereas the potential energy is composed of the Josephson energies (cosine terms) and the energy stored in the inductances of the circuit. From now on, we assume the SQUID inductances \( L_1 = L_2 \) and introduce the inductance asymmetry \( \eta \) defined by \( L_3 = \eta L_S \) [23]. Although the Lagrangian depends on six variables, the effective low-energy behavior of the circuit is determined by three variables only, as we will see below.
Let us first consider the phase variable $\gamma$. Its dynamics is that of a harmonic oscillator of frequency $\omega_\gamma = 1/\sqrt{L_S C_\gamma}$. Here $C_\gamma$ is the capacitance of the $\gamma$ node and $L_S = \eta(1 - \eta)L_S/(1 + \eta)$. Using the circuit parameters of Ref. [11] (see Appendix B), where $C_\gamma$ is estimated to be smaller than 0.1 fF, the frequency $\omega_\gamma$ is estimated to be larger than 1 THz, i.e., larger than all the other frequencies of the circuit. Next consider the phase variable $\xi$. Again using the circuit parameters of Ref. [11], we find that the characteristic inductive currents $\phi_0 L_{1,2}$ are on the order of 1 $\mu$A, much larger than the Josephson critical current $I_c^x$ of about 1 nA. In other words, the SQUID inductance is much smaller than the Josephson inductance $\sim 1/I_c^x$. Therefore, in a first approximation, the dynamics of $\xi$ is also of the harmonic oscillator type with a frequency $\omega_\xi = 1/\sqrt{L_S C_T^\pi}$. We estimate $\omega_\xi$ to be around 640 GHz. This frequency is much smaller than $\omega_\gamma$, but still quite high compared to the frequencies characterizing the dynamics of the variables $\varphi_1$, $\varphi_2$, $\psi$ and $\theta$ (around 10 GHz, see below). This implies that we can use the adiabatic approximation to eliminate the fast variables $\gamma$ and $\xi$ and obtain an effective Lagrangian for the slow variables $\varphi_1$, $\varphi_2$, $\xi$, and $\theta$.

In order to implement the adiabatic approximation, we write $\gamma = \gamma_0 + \delta \gamma$ and similarly $\xi = \xi_0 + \delta \xi$. Here $\gamma_0$ and $\xi_0$ correspond to the minima of the harmonic potentials confining the motion of these variables,

$$\gamma_0(x, y) = x + \eta y + \frac{\phi_S}{2}(1 - \eta) + \frac{1}{4\phi_0} L_S I_p(1 - \eta^2),$$

$$\xi_0(x, y) = x - \eta y - \frac{\phi_S}{2}(1 - \eta) + \frac{1}{4\phi_0} L_S I_p(1 - \eta^2),$$

where $x = (\varphi_1 + \varphi_2)/2$ and $y = (\varphi_2 - \varphi_1)/2$. Note that $\gamma_0$ and $\xi_0$ depend on $x$ and $y$, thus they acquire slow dynamics through the phases $\varphi_1$ and $\varphi_2$. The dynamics of the deviations $\delta \gamma$ and $\delta \xi$ is fast, determined by the frequencies $\omega_\gamma$ and $\omega_\xi$, respectively. Substituting the above decomposition for $\gamma$ and $\xi$ in the Lagrangian $[2]$, $[3]$, and $[4]$, and averaging over the fast variables $\delta \gamma$ and $\delta \xi$ we find the effective low-frequency potential energy

$$V(x, y, \psi) = 2E_J \left( -\cos(x) \cos(y) - s(\eta y + x) + b(y - y_B)^2 \right) - E_{J1}^T \cos(\psi) - E_{J2}^T \cos(\psi - \xi_0(x, y) + \phi_T).$$

Here we defined the reduced parameters $b = \phi_0/(L_I 0)$, $s = I_b/(2I_0)$ and $y_B = \pi/(\Phi_0/\phi_0)$. Furthermore, $E_J = \phi_0 I_0$, $E_{J1}^T = \phi_0 I_1^x$ and $E_{J2}^T = \phi_0 I_2^x$ are the different Josephson energies of the circuit. Note that the fast oscillations of $\xi$ lead to a renormalization of $I_2^x$; the value it takes in the effective low-frequency potential $[7]$ will therefore be smaller than the bare value appearing in $[3]$. Assuming that the bias current $I_b$ and the flux $\Phi_S$ are constant, we deduce from Eqs. $[5]$ and $[6]$ that $\gamma_0 = \dot{x} + \eta \dot{y}$ and $\xi_0 = \dot{x} - \eta \dot{y}$. The kinetic part of the Lagrangian can then be expressed as:

$$T(\dot{x}, \dot{y}, \dot{\psi}, \dot{\theta}) = \phi_0^2 \left( C_0 \dot{x}^2 + C_0 \dot{y}^2 + \frac{1}{2} C_1^T (\dot{x} - \eta \dot{y} - \psi)^2 + \frac{1}{2} C_2^T \dot{\psi}^2 + \frac{1}{2} C_2^T (\dot{x} - \psi - \theta)^2 + \frac{1}{2} C_P \dot{\theta}^2 \right).$$

The four variables of the Lagrangian can be separated in three groups. Indeed, the dynamics of $x$ and $y$ corresponds to that of the dc-SQUID, whereas the dynamics of $\psi$ is associated with that of the ACPT. The variable $\theta$ is used to model the effect of the gate voltage (cf. Sec. [11]). Note that the last term of the potential $[7]$ and the third term of the kinetic term $[8]$ couple the variables of the dc-SQUID and those of the ACPT together, and therefore are responsible of the coupling between these two elements.

E. Choice of variables for the dc-SQUID

The Lagrangian of the circuit is a function of the variables $x$ and $y$ associated with the SQUID. We want to change these variables to more appropriate ones which will be used below to describe the dynamics of the dc-SQUID (cf. Sec. [11]). Let us introduce the two-dimensional potential of the dc-SQUID $V_S(x, y)$ which is the contribution to the potential $V(x, y, \psi)$, Eq. $[7]$, depending only on the variables $x$ and $y$. It reads:

$$V_S(x, y) = 2E_J \left( -\cos(x) \cos(y) - s(\eta y + x) + b(y - y_B)^2 \right).$$

We stress here that this potential is identical to the one of a dc-SQUID alone, as studied by J. Claudon et al. [28]. The dynamics of the dc-SQUID is similar to that of a fictitious particle of mass $\approx \phi_0^2 2C_0$ which evolves in the potential $V_S(x, y)$. This potential undulates due to the product of cosine terms and contains wells that are separated by saddle points (see Fig. 3); $V_S(x, y)$ is modulated by the bias current $I_b$ and the flux $\Phi_S$. We consider now the case that the particle is trapped in one of these wells associated with a given local minimum $(x_0, y_0)$. Let us introduce the displacement variables around $(x_0, y_0)$ defined by $X = x - x_0$ and $Y = y - y_0$. We assume that the particle’s motion does not extend far from $(x_0, y_0)$. Then, we can replace the potential $V_S(x, y)$ by its third order expansion around $(X = 0, Y = 0)$. This expansion contains a cross-term in $XY$ which disappears by performing a rotation of the $(X,Y)$ plane (Fig. 3b) by the angle $\theta$, where $\theta$ is given by $\tan(\theta) = 2 (\partial_x^2 U(x_0, y_0))/(\partial_y^2 U(x_0, y_0) - \partial_y^2 U(x_0, y_0))$. The new variables of the dc-SQUID $X_\parallel$ and $Y_\perp$ associated with the rotated plane are defined by

$$\begin{pmatrix} X_\parallel \\ Y_\perp \end{pmatrix} = \begin{pmatrix} \cos(\theta) & \sin(\theta) \\ -\sin(\theta) & \cos(\theta) \end{pmatrix} \begin{pmatrix} X \\ Y \end{pmatrix}$$

and correspond to the position of the particle along the longitudinal and the transverse direction, defined by the
two-dimensionality of the SQUID potential. When the inductance $L_S$ is zero, the dynamics of the SQUID is described by only one variable, $x$, since $y = y_B$. In that case, the prefactors $\alpha_\parallel$ and $\alpha_\perp$ are equal to one. The two last terms of the potential contain one variable of the SQUID and one of the transistor. They thus couple the SQUID and the transistor. Note that the coupling terms of the second order and beyond have been neglected in the potential. The kinetic term of the Lagrangian can be rewritten as, using the variables $X_\parallel$ and $Y_\perp$,

$$T(X_\parallel, Y_\perp, \psi, \theta) =$$

$$\frac{1}{2} \phi_0^2 \left[ (2C_0 + \alpha_\parallel^2 C_2^T) \dot{X}_\parallel^2 + (2C_0 + \alpha_\perp^2 C_2^T) \dot{Y}_\perp^2 + C_\Sigma \dot{\psi}^2 ight.$$  

$$-2C_\frac{2}{2} \left( \alpha_\parallel \dot{X}_\parallel - \alpha_\perp \dot{Y}_\perp \right) \psi - 2C_\frac{2}{2} \alpha_\parallel \alpha_\perp \dot{X}_\parallel \dot{Y}_\perp$$

$$-2C_g \dot{\psi} \dot{\theta} + (C_P + C_g) \dot{\theta}^2 \right],$$

(13)

where $C_\Sigma$ is the total capacitor of the transistor, defined by $C_\Sigma = C_1^T + C_2^T + C_g$. The final expression for the total Lagrangian is obtained from $\mathcal{L}(X_\parallel, Y_\perp, \psi, \theta, X_\parallel, Y_\perp, \psi) = T(X_\parallel, Y_\perp, \psi, \theta) - V(X_\parallel, Y_\perp, \psi)$; it will be used in the next section to establish the Hamiltonian of the circuit.

**F. The Hamiltonian**

The Hamiltonian of the coupled circuit is a function of the variables $(X_\parallel, Y_\perp, \psi, \theta)$ and the conjugate momenta $(-hP_\parallel, -hP_\perp, -hn, -hQ)$. These momenta are related to the velocities $(\dot{X}_\parallel, \dot{Y}_\perp, \dot{\psi}, \dot{\theta})$ by the well-known expressions $-hP_\parallel \equiv \partial \mathcal{L}/\partial \dot{X}_\parallel, -hP_\perp \equiv \partial \mathcal{L}/\partial \dot{Y}_\perp, -hn \equiv \partial \mathcal{L}/\partial \dot{\psi}$ and $-hQ \equiv \partial \mathcal{L}/\partial \dot{\theta}$. The analytical expressions for the conjugate momentum variables are given in appendix C. The conjugate variables generate the charges $P_\parallel$, $n$ and $Q$ with unit $[-2e]$. We stress that these charges have a clear physical meaning. Indeed, $P_\parallel$ corresponds to the number of Cooper pairs stored in the two capacitors $C_0$; $n$ is the number of Cooper pairs on the island. In Eq. (14), the charge $-2enQ$ is equal to the bias charge $Q_b$. Performing the limiting procedure $C_P, Q_b \to \infty$, keeping their ratio constant $Q_b/C_P = V_g$, we see that the velocity $\dot{\theta}$ is constant and defined by $\dot{\theta} = V_g/\phi_0$. The expression for the Hamiltonian is determined by the Legendre transformation [22]:

$$\mathcal{H} = -hP_\parallel \dot{X}_\parallel - hP_\perp \dot{Y}_\perp - hn \dot{\psi} - \mathcal{L}. $$

(14)

All the velocities which appear in the Lagrangian have to be replaced by the conjugate variables, inverting the set of equations given in appendix C. The full Hamiltonian
of the circuit then takes the form:

$$\hat{H}(\vec{P}_{||}, \vec{X}_{||}, \vec{P}_{\perp}, \vec{Y}_{\perp}, \hat{n}, \hat{\psi}) =$$

$$\frac{(2e)^2}{2C_{||}} \vec{P}_{||}^2 + \frac{1}{2} k_{||} \vec{X}_{||}^2 + \delta \vec{X}_{||}^2$$

$$+ \frac{(2e)^2}{2C_{\perp}} \vec{P}_{\perp}^2 + \frac{1}{2} k_{\perp} \vec{Y}_{\perp}^2 + \delta_{\perp} \vec{Y}_{\perp}^2$$

$$+ \frac{(2e)^2}{C_{n||}} \vec{P}_{||} \vec{P}_{\perp} + \beta_{||} \vec{Y}_{\perp} \vec{X}_{||} + \beta_{\perp} \vec{Y}_{\perp} \vec{X}_{||}^2$$

$$+ \frac{(2e)^2}{C_{n\perp}} \vec{P}_{\perp} \left( \hat{n} - \frac{C_{g} V_{g}}{2e} \right)^2 - E_{1j}^T \cos(\hat{\psi}) - E_{2s}^T \cos(\hat{\psi} - \delta)$$

$$+ \frac{(2e)^2}{C_{n\parallel}} \vec{P}_{\perp} \left( \hat{n} - \frac{C_{g} V_{g}}{2e} \right) - E_{T} \alpha_{||} \sin(\hat{\psi})$$

$$- \frac{(2e)^2}{C_{n\perp}} \vec{P}_{\perp} \left[ \hat{n} - \frac{C_{g} V_{g}}{2e} \right] + E_{2T} \alpha_{\perp} \vec{Y}_{\perp} \sin(\hat{\psi} - \delta), \quad (15)$$

where the analytical expressions of the capacitances $C_{||}$, $C_{\perp}$, $C_{T}$, $C_{n||}$, $C_{n\perp}$, and $C_{\perp \perp}$ are given in Tab. [1]. Applying the standard canonical quantization rules, the classical variables have been replaced by their corresponding quantum operators. For more clarity, the conjugate pairs appear in the Hamiltonian in different colors. They satisfy the following commutation operations

$$\begin{align*}
\{\vec{X}_{||}, \vec{P}_{||}\} &= -i \\
\{\vec{Y}_{\perp}, \vec{P}_{\perp}\} &= -i \\
\{\hat{\psi}, \hat{n}\} &= -i
\end{align*}$$

(16)

The properties of the Hamiltonian (15) are not trivial. It describes the quantum dynamics of three sub-systems: the longitudinal and transverse phase oscillations within the dc-SQUID and the charge dynamics of the ACPT. Moreover the Hamiltonian describes the dominant coupling between the different quantum sub-systems. Very complex dynamics can appear in this full circuit. In this paper we mainly concentrate on the dynamics of the longitudinal SQUID phase mode and the charge dynamics of the ACPT and as well as on their coupling. The next section is dedicated to the study of the dc-SQUID Hamiltonian. We will deduce the simplified Hamiltonian for the longitudinal phase mode and justify why transverse phase mode can be neglected in this study.

### III. DC-SQUID

The dc-SQUID potential has already been discussed in Sec. [11] where we introduced the change of the variables $x$ and $y$ to the variables $X_{||}$ and $Y_{\perp}$. In this section, we first analyze the properties of the dc-SQUID potential in more detail. Then, we study the quantum dynamics of the dc-SQUID which is equivalent to that of a fictitious particle trapped in one of the wells of the potential. We will see under which conditions the dc-SQUID behaves as a phase qubit. In this Section, the coupling between the SQUID and the ACPT will be ignored such that we can consider the dc-SQUID as an independent element.

#### A. dc-SQUID potential

For an appropriate choice of bias parameters, i.e., the bias current $I_0$ and the flux $\Phi_S$, the SQUID potential contains wells that can be regrouped in families $[f]$ [24]. The index $f$ for the family $[f]$ is related to the number $f$ of flux quanta trapped in the dc-SQUID loop. The wells of the same family are located along the direction $x$, periodically spaced by a distance $2\pi$, and separated by saddle points. Since these wells have exactly the same geometry, the physical properties of the SQUID are independent of the particular well in which the fictitious particle is localized. Wells of the family $[f]$ exist only if the bias current satisfies the relation $I_c^f < I_0 < I_{c,\pm}^f$, where $I_c^f$ and $I_{c,\pm}^f$ are the positive and negative critical current of the family $[f]$, respectively. When $I_0 = I_{c,\pm}^f$, the local minima of the wells of the family $[f]$ and their closest saddle points coincide. The critical currents depend strongly on the flux $\Phi_S$ as shown in Fig. [4] for the experimental parameters of the circuit studied in Ref. [1]. Here, for a given flux, the absolute values of the critical currents $I_{c,\pm}^f$ and $I_{c}^f$ are different. This difference originates from the finite inductance asymmetry of the dc-SQUID ($\eta = 0.28$) and disappears when $\eta = 0$. For almost any value of the flux, the potential is characterized by a unique family of wells, except in the region close to $\pm\Phi_0/2$ where two families can coexist. This specific region has been investigated recently in order to study the double escape path of the particle, as well as to make the SQUID insensitive in first order to current fluctuations [14]. In the following we will discuss the dynamics of the dc-SQUID for the case where the fictitious particle is trapped in one of the wells of the potential.
Table I. Analytical expressions for the capacitances present in Hamiltonian 15 and the approximated expressions in the limit of $C_4 < C_2^1$ & $C_2^1 < C_0$. We have used the definitions $C^2 = 2C_0(C_2 + C_0(1 + \eta^2))$, $C_2 = C_2^1 + C_2^2$, and $C_R = C_0(C_2^1 + C_2^2)/2C_0$. The numerical values have been calculated for the parameters of the circuit studied in Ref. 1: $C_0 = 227 \text{ fF}$, $C_2 = 2.0 \text{ fF}$, $C_2^2 = 0.9 \text{ fF}$, $C_g = 29 \text{ aF}$, and for a zero escape angle ($\theta = 0$).

B. Hamiltonian of the dc-SQUID

In the full Hamiltonian 15 of the coupled circuit, we isolate the terms which only contain the operators $\hat{X}_\|$, and $\hat{Y}_\perp$ and their conjugate momenta $\hat{P}_\|$, and $\hat{P}_\perp$. These terms constitute the Hamiltonian of the dc-SQUID which takes the form:

$$
\hat{H}_S = E_C^\| \hat{P}_\|^2 + \frac{1}{2} k_{\|} \hat{X}_\|^2 + \alpha_{\|} \hat{X}_\| ^3
+ E_C^\perp \hat{P}_\perp^2 + \frac{1}{2} k_{\perp} \hat{Y}_\perp ^2 + \alpha_{\perp} \hat{Y}_\perp ^3
+ \beta_a \hat{Y}_\perp \hat{X}_\| + \beta_b \hat{Y}_\perp \hat{X}_\perp + \frac{(2e)^2}{C_{\|}} \hat{P}_\| \hat{P}_\perp, \tag{17}
$$

with the charging energies $E_C^\| = (2e)^2/(2C_0)$ and $E_C^\perp = (2e)^2/(2C_{\perp})$. The first three terms correspond to the Hamiltonian of a fictitious particle of mass $m = \phi_0^2 C_0$ which is trapped in an anharmonic potential along the longitudinal direction $X_\|$. These terms describe the dynamics of an anharmonic oscillator of characteristic frequency $\nu_p = \sqrt{k_{\|}/m_{\|}}/(2\pi)$. The next three terms correspond to the Hamiltonian of a fictitious particle of mass $m_{\perp} = \phi_0^2 C_{\perp}$ which is trapped in an anharmonic potential along the orthogonal direction $X_{\perp}$. These terms describe the dynamics of an anharmonic oscillator of characteristic frequency $\nu_{\perp} = \sqrt{k_{\perp}/m_{\perp}}/(2\pi)$. The anharmonicity of the two oscillators is due to the cubic term which results from the non-linearity of the Josephson junction. The last three terms of the dc-SQUID Hamiltonian mix the operators of the two oscillators and consequently couple them. Finally, the two-dimensional dynamics of the SQUID is similar to the dynamics of two coupled, one-dimensional oscillators. We proceed by introducing the following dimensionless operators: $\hat{X}_\| = \sqrt{\hbar \nu_p/2E_C^\|} \hat{X}_\|$, $\hat{P}_\| = -\sqrt{2E_C^\|/\hbar \nu_p} \hat{P}_\|$, $\hat{Y}_\perp = \sqrt{\hbar \nu_{\perp}/2E_C^\perp} \hat{Y}_\perp$ and $\hat{P}_\perp = -\sqrt{2E_C^\perp/\hbar \nu_{\perp}} \hat{P}_\perp$, which verify the commutation relations $[\hat{X}_\|, \hat{P}_\|] = i$ and $[\hat{X}_\perp, \hat{P}_\perp] = i$. With these new operators, the dc-SQUID Hamiltonian can be rewritten as

$$
\hat{H}_S = \frac{1}{2} \hbar \nu_p (\hat{P}_\|^2 + \hat{X}_\|^2) - \sigma \hbar \nu_p \hat{X}_\|^3
+ \frac{1}{2} \hbar \nu_{\perp} (\hat{P}_\perp^2 + \hat{Y}_\perp^2) - \sigma_{\perp} \hbar \nu_{\perp} \hat{Y}_\perp^3
+ \hbar \nu_a \hat{Y}_\perp \hat{X}_\| + \hbar \nu_b \hat{Y}_\perp \hat{X}_\perp + \hbar \nu_c \hat{P}_\| \hat{P}_\perp, \tag{18}
$$

where the parameters $\sigma$ and $\sigma_{\perp}$ correspond to the relative amplitude of the cubic term compared to the quadratic term and hence are direct measure of the degree of anharmonicity of the oscillators. The energies $\hbar \nu_a$, $\hbar \nu_b$ and $\hbar \nu_c$ are the coupling energies between the two oscillators.

Hereafter, we suppose that the particle is trapped in one of the wells of the family $[0]$. The geometry of this well varies as a function of the bias point ($I_0, \Phi_S$) of the circuit but does not change with the gate voltage $V_g$ and the flux $\Phi_T$. Therefore, the different parameters of the dc-SQUID Hamiltonian only depend on $I_0$ and $\Phi_S$. The numerical calculations of the parameters of the Hamiltonian under various biasing conditions, using the experimental parameters of the circuit studied in Ref. 1 (see Appendix B). Figs. 5(a) and 5(b) show the dependence of the transverse ($\nu_p$) and orthogonal ($\nu_{\perp}$) frequencies as a function of the bias point. Generally, the frequency $\nu_{\perp}$ is always higher than 35 GHz and is at least twice higher than $\nu_p$. The frequency $\nu_p$ tends towards zero when the bias point approaches the critical current line. Figs. 5(c,d,e) show the dependence of the parameters of the Hamiltonian as a function of the flux $\Phi_S$ for a fixed bias current of 1.89 $\mu$A. Fig. 5(d) shows the two anharmonicity parameters of the two oscillators. The anharmonicity $\sigma$ is typically around 3 % and increases close to the critical current line. The parameter $\sigma_{\perp}$ is very small regardless of the bias, which leads to us to the conclusion that the orthogonal oscillator can be considered as a harmonic one. Fig. 5(e) shows the different coupling energies. The coupling frequency $\nu_{\perp}^*$ of the order of 10 MHz and depends only weakly on $\Phi_S$. The coupling frequencies $\nu_a^*$ and $\nu_{\perp}^*$ are generally much higher than $\nu^*$ and depend on $\Phi_S$. Note that $\nu_{\perp}^*$ vanishes close to $\Phi_S = 0$ and $\nu_{\perp}^*$ is always negative. Numerical values of the parameters of the Hamiltonian, for $\Phi_S = 0.1 \Phi_0$ and $I_0 = 1.89 \mu$A, are given in Tab. II. The potential barrier (Fig. 5(f)) corresponds to the energy which separates the local minimum of the well from its closest saddle point. This energy is not a parameter of
FIG. 5. Numerical calculations of the parameters of the dc-SQUID Hamiltonian (17), using the experimental parameters of the circuit studied in Ref. [1]. The fictitious particle associated to the SQUID is trapped in a potential well of the family [0].

(a) Frequencies of the transverse and (b) orthogonal oscillators, i.e. $\nu_p$ and $\nu_{\perp}$, respectively, as a function of the flux $\Phi_S$ and the bias current $I_b$. The black line shows the measured switching current of the dc-SQUID from Ref. [25] which is close to the critical current. (c) Frequencies $\nu_p$ and $\nu_{\perp}$ as a function of $\Phi_S$ for $I_b = 2$ $\mu$A, i.e. for the bias points located on the dashed line of the Figs. (a) and (b). Measured frequencies $\nu_S$ from Ref. [25] are shown as red points. (d) Anharmonicity parameters $\sigma$ and $\sigma_{\perp}$ of the longitudinal and orthogonal oscillators respectively as a function of $\Phi_S$ for $I_b = 2$ $\mu$A. (e) Different coupling frequencies $\nu^*_a$, $\nu^*_b$ and $\nu^*_c$ between the two oscillators associated with the terms in $\hat{Y}_{\perp}\hat{X}_{\parallel}$, $\hat{Y}_{\parallel}\hat{X}_{\perp}$ and $\hat{P}_{\perp}\hat{P}_{\parallel}$ of $\hat{H}_S$ (17). (f) Energy $\Delta U$ for the potential barrier as a function of $\phi_S$ for $I_b = 2$ $\mu$A.

| Parameter | Value |
|----------|-------|
| $\nu_p$ | 16.24 GHz |
| $\sigma$ | 3.4 % |
| $\nu_{\perp}$ | 38.26 GHz |
| $\sigma_{\perp}$ | -0.006 % |
| $\nu^*_a$ | -612 MHz |
| $\nu^*_b$ | -348 MHz |
| $\nu^*_c$ | 12 MHz |

TABLE II. Numerical values of the parameters of the dc-SQUID Hamiltonian at the working point $\Phi_S = 0.1 \Phi_0$ and $I_b = 1.89 \mu$A, calculated using the experimental parameters of the circuit studied in Ref. [1].

the dc-SQUID Hamiltonian [18]. Nevertheless, if $\Delta U$ is sufficiently small, typically on the order of $\nu_p$, the expression [18] of $\hat{H}_S$ is too simplified. This occurs when the bias point is close to the critical current line. One should then take into account the coupling of the quantum levels inside the well to those outside the well. Note that this coupling is responsible of the escape of the particle from the well [26, 27]. This coupling will be neglected in the following, assuming that the particle is always trapped in a sufficiently deep well.
As \( \nu_L \gg \nu_B \), the quantum dynamics of the transverse oscillator is much faster than that of the longitudinal oscillator. We will assume in the following that the transverse oscillator is always in its ground state. It allows us to replace in \( \hat{H}_S \) the operators of the transverse oscillator by their average values given by \( \langle \hat{Y}_L \rangle = 0 \), \( \langle \hat{Y}_L^1 \rangle = 1/2 \) and \( \langle \hat{P}_L \rangle = 0 \). Only one of the three coupling terms remains after this simplification. The coupling reads \( h
u_p\hat{X}_L/2 \) and can be seen as a modification of the bias current \( I_B \) of less than 0.5 nA. This term will be neglected in the following. Under this condition, the dynamics of the particle along the longitudinal direction is given by the Hamiltonian

\[
\hat{H}_S = \frac{1}{2}h\nu_p (\hat{P}_L^2 + \hat{X}_L^2) - \sigma h\nu_p \hat{X}_L^3
\]  

In the following, the SQUID dynamics will be studied, using this simplified Hamiltonian. We denote by \( |n\rangle \) and \( E_n \) the eigenstates and the associated eigenenergies of \( \hat{H}_S \), respectively, such that \( \hat{H}_S |n\rangle = E_n |n\rangle \), where \( n \) is an integer number larger or equal to zero. If the anharmonicity is weak (\( \sigma \ll 1 \)), the energies \( E_n \) are given by a straightforward perturbative calculation and we find \( E_n = (n + 1/2)h\nu_p - 15/4\sigma^2(n + 1/2)^2h\nu_p \). Fig. 5 shows the approximate potential of the dc-SQUID and the three first eigenenergies. When the anharmonicity of the dc-SQUID is sufficiently large, the dynamics of the dc-SQUID in the presence of an external microwave perturbation involves only the two first levels \( |0\rangle \) and \( |1\rangle \). In that case, the dc-SQUID behaves as a qubit. Since for the dc-SQUID the Josephson energy is much larger than the charging energy, the fluctuations of the phase \( \hat{X}_L \) are much smaller than those of the charge \( \hat{P}_L \). For this reason, the dc-SQUID is referred to as a phase qubit. The dc-SQUID Hamiltonian can be rewritten in the basis \( \{|i\rangle \} \), using the Pauli matrices (See Appendix D), as

\[
\hat{H}_S = h\nu_p/2\sigma^2 \psi \equiv (E_i - E_0)/h \]  

is the characteristic qubit frequency. The frequency \( \nu_S \) can be approximated in first order with respect to the anharmonicity \( \sigma \) as \( h\nu_S = h\nu_p(1 - (15/2)\sigma^2/2) \). We see that the frequency \( \nu_S \) equals the plasma frequency \( \nu_p \) if the dc-SQUID anharmonicity is zero and decreases with increasing anharmonicity (see Fig. 5(c)). When the anharmonicity \( \sigma \) is zero, the dc-SQUID behaves as a harmonic oscillator described by the typical Hamiltonian \( \hat{H}_S = h\nu_p(\hat{a}\hat{a}^\dagger + 1/2) \), where \( \hat{a}^\dagger \) and \( \hat{a} \) are the one-plasmon creation and annihilation operators, respectively.

### IV. ASYMMETRIC COOPER PAIR TRANSISTOR (ACPT)

The dc-SQUID Hamiltonian is only one of the terms that appear in the full Hamiltonian of the circuit. It is coupled to a second term associated with the asymmetric Cooper Pair transistor (ACTP). This section is dedicated to the theoretical analysis of the ACPT dynamics, neglecting its coupling with the dc-SQUID. We will first build and analyze the ACPT Hamiltonian and will show that the ACPT can be viewed as a charge qubit. We will then describe the ACPT by using only two charge states. The errors on the level of the eigenenergies and the eigenstates induced by this simplified description will be estimated.

A. ACPT Hamiltonian

The ACPT Hamiltonian is identified by isolating from the full Hamiltonian of the circuit the terms which only contain the operators \( \hat{\psi} \) and \( \hat{n} \). After some straightforward algebra, the ACPT Hamiltonian reads

\[
\hat{H}_{ACPT} = E_C^T (\hat{n} - n_g)^2 - E_J^T \cos(\delta/2) \cos(\hat{\psi}/2) + \mu \sin(\delta/2) \sin(\hat{\psi}/2),
\]

where \( E_J^T \equiv E_{J_1}^T - E_{J_2}^T \) and \( E_C^T \equiv (2e)^2/(2C_T) \) are the Josephson and charging energies of the ACPT, respectively. It is the generalization of the Quantronium Hamiltonian \[ H_{ACPT} \] for which asymmetries in critical current and capacitance of the Cooper pair transistor were neglected. The gate-charge \( n_g \equiv C_g V_g/(2e) \) corresponds to the number of Cooper pairs induced by the voltage applied to the gate-capacitance. The charge states \( |n_{2e}\rangle \) are the eigenvectors of the charge operator \( \hat{n} \), such \( \hat{n}|n_{2e}\rangle = n|n_{2e}\rangle \) where \( n = 0, \pm 1, \pm 2... \) is the number of excess Cooper pairs on the transistor island. Using the commutation relation \( [\hat{\psi}, \hat{n}] = i \), we identify the action of the operator \( e^{i\hat{\psi}} = \sum_n |n_{2e}\rangle \langle (n + 1)_{2e}| \) which decreases the number of Cooper pairs on the island by one unit. In the charge representation, the ACPT Hamiltonian can be written as

\[
\hat{H}_{ACPT} = E_C^T (\hat{n} - n_g)^2 - \rho_J(\delta/2)/2 \left[ \sum_n e^{-i(\delta/2 + \chi)}|n_{2e}\rangle \langle (n + 1)_{2e}| \right.
\]

\[
\left. + \sum_n e^{i(\delta/2 + \chi)}|n_{2e}\rangle \langle n_{2e}| \right],
\]

with \( E_J^T/2(\cos(\delta/2) + i\mu \sin(\delta/2)) \equiv \rho_J(\delta/2)2e^{i\chi} \), \( \tan \chi = \mu \tan(\delta/2) \) and \( \rho_J(\delta) = E_J^T/2(\cos^2(\delta/2) + \mu^2 \sin^2(\delta/2)) \).
The ACPT Hamiltonian is composed of a charging and a Josephson term which are proportional to \( E_C^T \) and \( \rho_j(\delta) \), respectively. Let us focus first on the case of a zero Josephson coupling (\( \rho_j = 0 \)). The eigenstates of the ACPT Hamiltonian are then the charge states \(|n_{2e}\rangle\) with the associated eigenenergies \( E_C^T(n - n_g)^2 \). Fig. 7 shows the energy spectrum of the ACPT as a function of \( n_g \) for a charging energy \( E_C^T = 1.28 \) K. This spectrum consists of a series of parabolas, each parabola being associated with a specific charge state \(|n_{2e}\rangle\) with a minimum energy for \( n_g = n \). Notice that when the energy difference between the ground charge state and the first excited charge state is much larger than \( k_B T \), the charge on the island is well quantized leading to the Coulomb blockade phenomena. For \( n_g = 0.5 \), the energy parabolas of the states \(|0_{2e}\rangle\) and \(|1_{2e}\rangle\) cross each other and the states \(|0_{2e}\rangle\) and \(|1_{2e}\rangle\) are degenerate. This degeneracy is lifted by the Josephson term which couples the neighboring charge states to each other. The amplitude of this Josephson coupling is given by \( \rho_j \). Fig. 8 shows the dependence of \( \rho_j \) on \( \delta \) for three different Josephson asymmetries. Since \( \rho_j \) is 2\( \pi \) periodic in \( \delta \), the range of \( \delta \) has been restricted to the interval between \(-2\pi\) and \(2\pi\). In the case of a symmetric transistor (\( \mu = 0 \)), the Josephson coupling is maximum for \( \delta = 0 \), equal to \( E_J^T \); it is zero at \( \delta = \pm \pi \). For a finite asymmetry \( \mu \), the Josephson coupling reaches a maximum for \( \delta = 0 \) and equals \( E_J^T \), and a minimum equals to \( \mu E_J^T \) for \( \delta = \pm \pi \). For a Cooper pair box (\( \mu = \pm 100\% \)), the Josephson coupling does not depend on \( \delta \) and remains equal to \( E_J^T \). The Josephson coupling then depends strongly on the Josephson asymmetry, especially for \( \delta = \pm \pi \) where it can vary from zero to \( E_J^T \). The full energy spectrum of the ACPT, which takes into account the Josephson coupling, is calculated numerically by diagonalizing the Hamiltonian in a basis of 8 charge states. It is plotted as a function of \( n_g \) in Fig. 7.

For a fixed superconducting phase \( \delta = 0 \) and for the parameters of [1]: \( E_C^T = 1.28 K \), \( E_J^T = 1.01 k_B K \) and \( \mu = -41.6\% \). The spectrum consists of several energy bands which do not cross each other, leading to an energy gap between the lowest and the first bands. We associate with these two bands the ground state \(|-\rangle\) and the first excited state \(|+\rangle\), respectively. These two states correspond to the states of a qubit, with a characteristic energy \( \hbar \nu_T \). Indeed, when the ACPT Hamiltonian is subject to an adequate perturbation, the quantum dynamics of the ACPT will only involve the states \(|-\rangle\) and \(|+\rangle\). Because \( E_J^T > E_C^T \) the states \(|-\rangle\) and \(|+\rangle\) can be written as a superposition of a few charge states (typically four). For this reason, the qubit formed by the ACPT is referred to as a charge qubit. The Hamiltonian of the charge qubit can be written using the Pauli matrix (see appendix D) as

\[
\frac{\hbar \nu_T}{2} (\sigma_z - T) \cdot \delta. \tag{22}
\]

The frequency \( \nu_T \) of the charge qubit depends on the two parameters \( \delta \) and \( n_g \) as shown in Fig. 8(a). This frequency, as is \( \rho_j \), is also 2\( \pi \) periodic as a function of \( \delta \). It is maximum and minimum at the points \( (\delta = 0, n_g = 1/2) \) and \( (\delta = \pi, n_g = 1/2) \), respectively. These are optimal points where the ACPT is, in first order, insensitive to charge, flux and current fluctuations. Fig. 10 shows the experimental frequency of the charge qubit measured in Ref. [1] as a function of \( \delta \). The theoretical fit, shown in red, is obtained by diagonalizing the ACPT Hamiltonian in a basis of 8 charge states. It allows us to accurately find the parameters of the charge qubit, i.e. the Josephson energy \( E_J^T = 21.08 \) GHz, the charging energy \( E_C^T = 26.76 \) GHz and the Josephson asymmetry \( \mu = -41.6\% \). Nevertheless, the capacitance asymmetry can not be extracted from the fit since it does not enter in

FIG. 7. Energy spectrum of the ACPT as a function of \( n_g \) in the case of a zero Josephson coupling (dashed lines) and in the case of \( \rho_j(\delta = 0) = 21.08 \) GHz (full lines). These energy spectra have been calculated numerically, using the parameters from Ref. [1]: \( E_C^T = 21.08 \) GHz, \( E_J^T = 26.76 \) GHz and \( \mu = -41.6\% \). The states \(|-\rangle\) and \(|+\rangle\) associated to the two lowest energy bands correspond to the states of the charge qubit.

FIG. 8. Josephson coupling \( \rho_j \) as a function of the superconducting phase \( \delta \) across the ACPT, for three different Josephson asymmetries: \( \mu = 0 \) for a symmetric transistor, \( \mu = -41.5\% \) for an asymmetric transistor and \( \mu = 100\% \) for a Cooper pair box.
with the energies $E_i$ for $i = 0, 1$. The eigenvalues of this simplified Hamiltonian $H_{ACPT}$ which reads, in the charge basis ($|n\rangle$),

$$H_{ACPT}^0 = \left( -\rho_j e^{i(\delta/2+\chi)/2} E_C^0 (1-n_g)^2 \right).$$

The eigenvalues of this simplified Hamiltonian are given by $E_{\pm}^0 = \frac{1}{2} E_C^0 \left( n_g^2 + (1-n_g)^2 \right) \pm \frac{1}{2} \sqrt{E_C^0 (1-2n_g)^2 + \rho_j^2}$, and the qubit energy reads $\hbar \nu_T^0 = E_+^0 - E_-^0$. Note that for $n_g = 0.5$, we have $\hbar \nu_T^0 = \rho_j$. The eigenstates $|\pm\rangle$ and $|\pm\rangle$, associated with the energies $E_-$ and $E_+$, can be written as a function of the charge states as

$$\begin{align*}
|+\rangle &= \alpha^*|02e\rangle + \beta|12e\rangle \\
|-\rangle &= -\beta^*|02e\rangle + \alpha|12e\rangle
\end{align*}$$

with $\alpha = \cos(\theta/2) e^{i(\delta/2+\chi)/2}$, $\beta = \sin(\theta/2) e^{i(\delta/2+\chi)/2}$ and $\tan \theta = \frac{E_C^0 (1-2n_g)}{\rho_j}$ ($\theta \in [0, \pi]$).

For the ACPT studied experimentally in Ref. [1], the condition $\rho_j(\delta) \ll E_C^0$ is too strong, especially at $\delta = 0$ where the ratio $\rho_j/E_C^0 \approx 79\%$ is maximum. For this reason, we would like to quantify the error induced by the description of the ACPT with only two charge states. Let us first focus on the error in the qubit energy. Fig. 11 shows the difference between the frequencies $\nu_T^0$ calculated with two charge states and the real frequency $\nu_T$ as a function of $n_g = 0.5$. This difference is minimum for $\delta = \pm \pi$ and equals 58 MHz, corresponding to an error in energy of 0.5%. For $\delta = 0$, the difference is maximum and equals 780 MHz, corresponding to an error in energy of 3.8%.

We now discuss the error made on the level of the states $|0\rangle$ and $|1\rangle$. The non-reduced Hamiltonian of the ACPT can be rewritten as a function of $H_{ACPT}^0$ as $H_{ACPT} = H_{ACPT}^0 + \tilde{W}$, where $\tilde{W}$ is a perturbative term which takes the form $\tilde{W} = -\frac{\rho_j(\delta)}{2} \sum_{n \neq 0 \neq 1} \left( e^{-i(\delta/2+\chi)}|n\rangle\langle n+1| + e^{i(\delta/2+\chi)}|n\rangle\langle n+1| \right) + E_C^0 \sum_{n \neq 0 \neq 1} (n-n_g)^2 |n\rangle\langle n|).$ Using first-order perturbation theory, we calculate the probability $|\langle -| \rangle|^2 (|\langle +| +\rangle|^2)$ that the state $|\rangle$ ($|+\rangle$) is in the state $|-\rangle$ ($|+\rangle$). Fig. 11 shows the dependence of these probabilities as a function of $\delta$ for $n_g = 0.5$. The probability $|\langle -| \rangle|^2$ is always smaller than the probability $|\langle +| +\rangle|^2$. Indeed, the energy of the state $|+\rangle$ is closer to the charg-
we consider in the following that

|−⟩calculation of the analytical expression of the coupling

closer-in-energy charge states. In order to simplify the

calculate the real frequency to |ν| and |−⟩numerical one. The probabilities

dependence of the probabilities |⟨−⟩| and |−⟩|⟨−⟩| decrease with increasing |E|

coulomb and |⟨−⟩| mixing more and decrease with increasing |E|

coupling. However, in the studied circuit the dc-SQUID

coupling Hamiltonian by considering the ACPT as a

tunable harmonic oscillator or a phase qubit. We will see that the

coupling is the sum of two distinct contributions: a

capacitive and an inductive Josephson coupling.

V. COUPLING

So far, we have considered independently the quantum

dynamics of the longitudinal mode of the dc-SQUID and

the ACPT. However, in the studied circuit the dc-SQUID

and the ACPT are connected in parallel and therefore
coupled to each other. The independent dynamics of

the ACPT and the dc-SQUID has to be reconsidered
especially when the two qubits are close to resonance

(νS ≈ νT). In this case the coupling effects are the

strongest. In this section, we derive the expression of

the coupling Hamiltonian by considering the ACPT as a

charge qubit and the dc-SQUID as either a tunable har-
monic oscillator or a phase qubit. We will see that the

total coupling is the sum of two distinct contributions: a

capacitive and an inductive Josephson coupling.

A. Capacitive coupling

The capacitive coupling Hamiltonian couples by definition

the charge of the dc-SQUID with that of the ACPT. It reads

\[ \hat{H}_{\text{Coupl,Capa}} = -\frac{(2e)^2}{C_n|V|} \sqrt{\frac{\hbar \nu_p}{2E_C^2}} \hat{P}(\hat{n} - n_g). \] (25)

We consider hereafter two different limits for the dc-

SQUID in order to simplify this capacitive coupling

Hamiltonian.

The first limit corresponds to a dc-SQUID with an

anharmonicity factor σ equal to zero. This limit can be

achieved when the fictitious particle associated with

the dc-SQUID is trapped in a deep well, which is gener-

ally true when the dc-SQUID is biased at the working

point \( I_b = 0 \) and \( \Phi_S = 0 \). Under this condition, the

\[]
dc-SQUID behaves as a harmonic oscillator and several levels are involved in the dynamics. The momentum operator in the charge coupling Hamiltonian is then given by \( \hat{P}_c = (\hat{a}^\dagger + \hat{a})/\sqrt{2} \), where \( \hat{a}^\dagger \) and \( \hat{a} \) are the one-plasmon creation and annihilation operators. If we describe the charge qubit with two charge states (Eq. 24), the operator \( \hat{n} \) can be written as \( \hat{n} = \frac{1}{2} \left[ 1 + \frac{i}{2} \sin(\theta) \hat{\sigma}_x^T - \frac{1}{2} \cos(\theta) \hat{\sigma}_y^T \right] \), where \( \hat{I} \) is the identity operator and \( \hat{\sigma}_x^T \) and \( \hat{\sigma}_y^T \) are the Pauli matrices defined in the eigenstate basis of the charge qubit. For \( n_g = 1/2 \), we have \( \theta = \pi/2 \) and the charge coupling Hamiltonian simplifies to

\[
\hat{H}_{\text{Coupl,Capa}} = -E_{c,c} (\hat{a}^\dagger + \hat{a}) \hat{\sigma}_x^T / \sqrt{2},
\]

where \( E_{c,c} = e^2 / C_n \parallel \sqrt{\hbar \nu_p / E_C^c} \) determines the strength of the capacitive coupling.

The second limit is realized when the anharmonicity \( \sigma \) is typically around 3 % or larger; then the dc-SQUID can be described as a phase qubit. Generally, \( \sigma \) does not exceed 10% (see Fig. 5(d)) so the two lowest eigenstates of the dc-SQUID Hamiltonian are very close to the two lowest eigenstates of the harmonic oscillator. Consequently, the momentum operator can be expressed in terms of the eigenbasis of the charge qubit as \( \hat{P}_c \approx (|0 \rangle \langle 0| + |1 \rangle \langle 1|) / \sqrt{2} = \hat{\sigma}_z^T / \sqrt{2} \). For \( n_g = 1/2 \), the charge coupling Hamiltonian takes the following form:

\[
\hat{H}_{\text{Coupl,Capa}} = -E_{c,c} \hat{\sigma}_z^T \hat{\sigma}_x^T.
\]

The capacitive coupling produces a transverse interaction between the two quantum systems. This coupling was already discussed in Ref. [30, 31] where a Cooper pair box was coupled to a harmonic oscillator. The capacitive coupling is vanishing in the limit of small \( C_T^c \). We notice also that \( E_{c,c} \) depends on the bias variables through \( \sqrt{\nu_p} \) and therefore is weakly tunable.

### B. Josephson coupling

The Josephson coupling Hamiltonian couples the phases \( \hat{X}_\parallel \) and \( \hat{\psi} \) associated with the dc-SQUID and the transistor, respectively. This coupling is mediated via the Josephson junction 2 of the ACPT and reads

\[
\hat{H}_{\text{Coupl,Jos}} = -E_{c,J} \cos(\delta/2 - \chi) \hat{\sigma}_y^T \hat{\sigma}_y^T - E_{c,J} \sin(\delta/2 - \chi) \hat{\sigma}_y^T \hat{\sigma}_y^T.
\]

The Josephson coupling contains two different terms. One describes a transverse coupling \( \hat{\sigma}_y^T \hat{\sigma}_y^T \) or \( (\hat{a} - \hat{a}^\dagger) \hat{\sigma}_y^T \) which gives rise to coherent energy exchange at the resonance between the charge qubit and the phase qubit (or the oscillator). The effects of this first term on the quantum dynamics of the circuit are similar to those produced by the capacitive coupling term in \( \hat{\sigma}_z^T \hat{\sigma}_x^T \) or \( (\hat{a} - \hat{a}^\dagger) \hat{\sigma}_x^T \). The second term \( \hat{\sigma}_y^T \hat{\sigma}_y^T \) or \( (\hat{a} - \hat{a}^\dagger) \hat{\sigma}_y^T \) contains a transverse contribution for the SQUID and a longitudinal term which depends on the transistor qubit state. Its contribution will explain the quantum measurement of the charge qubit in the \( \nu_S \ll \nu_T \) limit (see Sec. IV C). We notice finally that the two terms are strongly tunable with the bias parameter \( \delta \).

### VI. QUANTUM DYNAMICS OF THE COUPLED CIRCUIT

The full Hamiltonian of the coupled circuit is given by the sum of the Hamiltonians of the dc-SQUID, the ACPT and the coupling. It reads: \( \hat{H} = \hat{H}_S + \hat{H}_{\text{ACPT}} + \hat{H}_{\text{Coupl,Capa}} + \hat{H}_{\text{Coupl,Jos}} \). In order to simplify this Hamiltonian, we consider hereafter the situation where the gate-charge is fixed to \( n_g = 0.5 \). This charge value has been mainly used for the charge qubits experiments (see Refs. [1, 16, 32]). Indeed, at \( n_g = 0.5 \), the charge qubit is insensitive in first order to charge noise and, therefore, its coherence time is longer. We will first derive the expression of the full Hamiltonian when the two quantum systems are close to the resonance condition. We will
consider the dc-SQUID either as a phase qubit or as a harmonic oscillator. Finally, we will discuss the quantum measurement scenario of the transistor states. Its principle derives from the coupling to the SQUID, enabling to use it as a detector of the charge qubit’s state.

A. A charge qubit coupled to a phase qubit

We first consider the dc-SQUID as a phase qubit such that the full Hamiltonian governs the quantum dynamics of two-coupled qubits. Using the expressions of the Josephson and capacitive coupling in that limit, we find

\[ \hat{H} = \frac{h \nu_S}{2} \hat{\sigma}^S_z + \frac{h \nu_T}{2} \hat{\sigma}^T_z - E_{c,c} \hat{\sigma}^S_z \hat{\sigma}^T_z + E_{c,j} \cos(\delta/2 - \chi) \hat{\sigma}^S_y \hat{\sigma}^T_y - E_{c,j} \sin(\delta/2 - \chi) \hat{\sigma}^S_z \hat{\sigma}^T_z \]

Let us introduce the raising operators \( \hat{\sigma}^T_z \) and \( \hat{\sigma}^S_z \) and the lowering operators \( \hat{\sigma}^T_\pm \) and \( \hat{\sigma}^S_\pm \). These operators are defined by \( \hat{\sigma}^T_\pm = (\hat{\sigma}^T_z \pm i \hat{\sigma}^T_y)/\sqrt{2} \) and \( \hat{\sigma}^S_\pm = (\hat{\sigma}^S_z \pm i \hat{\sigma}^S_y)/\sqrt{2} \).

The terms \( \hat{\sigma}^S_z \hat{\sigma}^T_z \) and \( \hat{\sigma}^S_\pm \hat{\sigma}^T_\pm \) can be written using the four products \( \hat{\sigma}^T_z \hat{\sigma}^T_\pm, \hat{\sigma}^S_\pm \hat{\sigma}^T_\pm, \hat{\sigma}^S_\pm \hat{\sigma}^T_\pm \) and \( \hat{\sigma}^S_\pm \hat{\sigma}^T_z \), whereas the term \( \hat{\sigma}^T_\pm \hat{\sigma}^S_\pm \) is a function of the operators \( \hat{\sigma}^S_z \hat{\sigma}^T_z \) and \( \hat{\sigma}^S_\pm \hat{\sigma}^T_\pm \).

The product \( \hat{\sigma}^S_\pm \hat{\sigma}^T_\pm \) corresponds to an excitation of the phase qubit and de-excitation of the charge qubit. This coupling term mediates the transition between the states \( |0, +\rangle \) and \( |1, -\rangle \). This transition is only relevant close to resonance, where it contributes to the low-frequency dynamics of the coupled system. Far away from resonance, this term gives rise to high-frequency dynamics (frequencies of the order of \( \nu_S \) and \( \nu_T \)) which is averaged away on the typical time scales of the experiment [11]. The rotating-wave approximation consists in neglecting these non-resonant terms, called also inelastic terms. The previous reasoning applied to the term \( \hat{\sigma}^T_\pm \hat{\sigma}^T_\pm \) leads to the same result. The products \( \hat{\sigma}^S_\pm \hat{\sigma}^T_\pm \) and \( \hat{\sigma}^S_\pm \hat{\sigma}^T_\pm \) couple the states \( |0, -\rangle \) and \( |1, +\rangle \). The transition between these two states leads to high-frequency dynamics and are neglected hereafter. For the same reason, the coupling term \( \hat{\sigma}^S_\pm \hat{\sigma}^T_\pm \) will be ignored. Finally, close to the resonance, the full Hamiltonian simplifies to

\[ \hat{H} = \frac{h \nu_S}{2} \hat{\sigma}^S_z + \frac{h \nu_T}{2} \hat{\sigma}^T_z - \frac{g}{2} (\hat{\sigma}^S_+ \hat{\sigma}^T_+ + \hat{\sigma}^S_- \hat{\sigma}^T_-), \quad (33) \]

where the coupling strength \( g \) is given by \( g = 2E_{c,c} - 2E_{c,j} \cos(\delta/2 - \chi) \). We assume hereafter that the gate capacitance \( C_g \ll C_T^2 \) and \( C_T^2 \ll C_0 \) which is true for the measured circuits in Refs. [11][12]. In these limits, we have \( C_{\parallel} \approx 2C_0(C_T^2 + C^2)/C_0(C_T^2) \). The coupling energy can then be written as a function of the capacitance (\( \lambda \)) and Josephson (\( \mu \)) asymmetries of the ACPT as:

\[ g = \frac{\alpha_{\parallel}}{2} \sqrt{\frac{E_{\parallel}^c C_0}{h \nu_p}} \left[ (1 + \lambda)h \nu_p - (1 + \mu)E_J^c \cos(\delta/2 - \chi) \right]. \quad (34) \]

Further simplification of the coupling energy can be obtained using the fact that the charge qubit is described with two charge states, and that the SQUID has a zero anharmonicity. Then, the transistor frequency is related to \( \delta \) by \( \nu_T = (E_J^c/h)\cos(\delta/2) + \mu^2 \sin^2(\delta/2) \) (see Sec. IV B) and we have \( \nu_p = \nu_S \). By using these relations, the coupling \( g \) (34) can be rewritten as

\[ g = \frac{\alpha_{\parallel}}{2} \sqrt{\frac{E_{\parallel}^c C_0}{h \nu_p}} \left[ (1 + \lambda)\nu_S - (1 + \mu)\nu_T (\cos^2(\chi) + \frac{\sin^2(\chi)}{\mu}) \right]. \quad (35) \]

Note that coupling \( g \) depends strongly on the bias parameters, via the phase \( \delta \) and the frequency \( \nu_S \). Therefore, the proposed circuit presents an intrinsic tunable coupling between a charge and a phase qubit, which we will analyze in more detail in Sec. VII. This transverse coupling enables to realize two-qubit gates operation as for example the \( \sqrt{\text{SWAP}} \) operation [33].

B. A charge-qubit coupled to a tunable harmonic oscillator

The Hamiltonian [33], which describes the dynamics of two coupled qubits, is valid if the second and highest levels of the dc-SQUID do not participate in the quantum dynamics of the circuit. This is the case when the anharmonicity of the SQUID is sufficiently strong. In the limit of zero anharmonicity, the quantum dynamics of the coupled circuit is described by the Jaynes-Cummings Hamiltonian [34]

\[ \hat{H} = \hbar \nu_p (\hat{a} \hat{a}^\dagger + \frac{1}{2}) + \frac{h \nu_T}{2} \hat{\sigma}^T_z - \frac{g}{2} (\hat{a}^\dagger \hat{\sigma}^T_+ + \hat{a} \hat{\sigma}^T_-). \quad (36) \]

This Hamiltonian (36) is very similar to the one obtained by coupling a charge [35] or a transmon [36] or phase qubit [37] to a coplanar waveguide cavity. In our circuit the resonator is realized by a micro dc-SQUID. It can be viewed as "micro-resonator" more convenient for integration than usual coplanar resonators since its size ranges in the micrometer scale three orders of magnitude smaller. Moreover its resonance frequency is strongly tunable. However up to now it suffers from a much shorter coherence time.

C. Quantum measurements of the charge qubit by the dc-SQUID

Our theoretical analysis introduces two different kinds of coupling which can be used to perform quantum measurements of the charge qubit by the dc-SQUID. Here we will apply this analysis to the quantronium read-out [16] and the adiabatic quantum transfer method used in Ref. [1].
1. Quantronium read-out

This read-out is obtained the limit where $\nu_T \gg \nu_S$. As we will see, even if the qubits are far from resonance the coupling still affects the dynamics of the circuit. Indeed, since the dynamics of the ACPT is much faster than the dc-SQUID dynamics, the dc-SQUID is only sensitive to the average value of the ACPT operators. After some straightforward algebra, the effective Hamiltonian of the dc-SQUID takes the following form:

$$\hat{H}_{S,\text{eff}}^{(\pm)} = \hat{H}_S - \frac{C_T}{C_n} \frac{\partial E_{\pm}}{\partial n_g} \tilde{P}_0 + \alpha \frac{\partial E_{\pm}}{\partial \delta} \vec{X},$$

with $E_{\pm}$ and $E_{\mp}$ the eigenenergies of the ACPT associated with the states $|+\rangle$ and $|\mp\rangle$. For $n_g = 0$ or $n_g = 1/2$, we have $\partial E_{\pm}/\partial n_g = 0$ and the SQUID Hamiltonian simplifies to

$$\hat{H}_{S,\text{eff}}^{(\pm)} = \hat{H}_S - \phi_0 I_{\text{add}}^{(\pm)} \vec{X},$$

where $I_{\text{add}}^{(\pm)} \equiv -\langle \alpha \rangle / \phi_0 \partial E_{\pm} / \partial \delta$ adds to the bias current $I_b$ and depends on the state of the ACPT $|\mp\rangle$ or $|+\rangle$. For $n_g = 1/2$, in the limit of two charge states only, the additional current reads

$$I_{\text{add}}^{(\pm)}(n_g = 1/2) = \pm \frac{E_T}{2\phi_0} (1 + \mu) \sin(\delta/2 - \chi).$$

The currents $I_{\text{add}}^{(+)}(n_g = 1/2)$ and $I_{\text{add}}^{(-)}(n_g = 1/2)$ take opposite values and, therefore, can be used to determine the transistor state. Fig. 13 shows the normalized current $I_{\text{add}}^{(-)}/(E_T/2\phi_0)$ for $n_g = 1/2$ as a function of the phase $\delta$ for an asymmetric ($\mu = 0.42$) and symmetric ($\mu = 0$) transistor and a Cooper pair box ($\mu = 1$). For a symmetric transistor, the current has a strong discontinuity at $\delta = \pm \pi$ jumping between the two extreme values $-E_T^2/2\phi_0$ and $E_T^2/2\phi_0$. This discontinuity disappears for a finite Josephson asymmetry; its maximum value is lower. Naturally, for a Cooper pair box, the additional current is zero as one of the junctions is replaced by a pure capacitance.

As the switching probability of the dc-SQUID depends strongly on the bias current [28] the additional current $I_{\text{add}}$ can be used to detect the ACPT state. For instance, in the Quantronium circuit [16], the state of a symmetric Cooper Pair Transistor (CPT) is read out by measuring the switching probability of a Josephson junction placed in parallel with the CPT and which replaces the dc-SQUID in the studied circuit. In that case, the charge qubit read-out is explained by the $\vec{X} \hat{\sigma}_z / \tilde{\sigma}_z$ term resulting from the Josephson coupling [29].

The expression [29] of the additional current has been established under the condition $\nu_T \gg \nu_S$. Is this expression still valid when $\nu_T \approx \nu_S$? In order to answer this question, we compare experimental data and theoretical predictions for the difference of the additional currents at $n_g = 1/2$ and $n_g = 0$ in the ground state of the ACPT. This difference noted $\Delta I_{\text{add}}^{(-)} \equiv I_{\text{add}}^{(-)}(n_g = 1/2) - I_{\text{add}}^{(-)}(n_g = 0)$ is shown in Fig. 14 as a function of the flux $\Phi_S$. The theoretical curve calculated without any free parameters agrees well with the experiment when $\nu_T / \nu_S \gg 1$, but differs from the experiment when $\nu_T \approx \nu_S$. The measured current amplitude drops when $\nu_T$ is close to $\nu_S$ suggesting a drop in the contrast of the quantum measurement when $\nu_T \approx \nu_S$. 

![FIG. 13. (Color online). Effective current $I_{\text{add}}^{(-)}$ at $n_g = 1/2$ as a function of the phase $\delta$ for an asymmetric ($\mu = 0.42$) and symmetric ($\mu = 0$) transistors and a Cooper pair box ($\mu = 1$).](image)

![FIG. 14. (Color online). (a) Experimental (black squares) and theoretical (red line) current difference $\Delta I_{\text{add}}^{(-)}$ at $n_g = 1/2$ as a function of the flux $\Phi_S$ (or the phase $\delta$ indicated in the upper scale). (b) Frequencies ratio $\nu_T/\nu_S$ as a function of $\Phi_S$.](image)
the state of the flux pulse, the state energy spectrum where an adiabatic transfer happens during spectroscopy (a). Inset shows an antilevel crossing in the nanosecond flux pulse which transfers the quantum state of spectroscopy (a). The coupling terms  \( \hat{\sigma}_+^T \hat{\sigma}_+^T \) and  \( \hat{\sigma}_-^T \hat{\sigma}_-^T \) produce an antilevel-crossing the amplitude of which depends on the coupling strength \( g \) (see Inset of Fig. 15(c)). During a nanosecond flux pulse, the coupled system remains in its original energy state; as a result, the quantum state \( |0, +\rangle \) evolves adiabatically into the state \( |1, -\rangle \). Due to large value of the coupling strength close to the escape point, Landau-Zener transitions can be neglected. The final state, i.e. \( |1, -\rangle \) or \( |0, -\rangle \), is determined by the switching measurement.

2. Adiabatic quantum transfer

In Fig. 15, the first energy levels of the charge and phase qubit in the coupled circuits are plotted. The quantum measurement of the charge qubit is performed by a nanosecond flux pulse which transfers the quantum state \( |0, +\rangle \) prepared at the working point \( (I_{wp}^{sc}, \Phi_{Swp}^{sc}) \) to the measurement point \( (I_{wp}^{sw}, \Phi_{Swp}^{sw}) \). At that point the SQUID is very close to the critical line and the escape probability is finite. Spectroscopy measurements show clearly the read-out of the charge qubit by this method even in the limit \( \nu_T \approx \nu_S \). How can we explain this read-out?

The coupling terms \( \hat{\sigma}_+^T \hat{\sigma}_+^T \) and  \( \hat{\sigma}_-^T \hat{\sigma}_-^T \) produce an antilevel-crossing the amplitude of which depends on the coupling strength \( g \) (see Inset of Fig. 15(c)). During a nanosecond flux pulse, the coupled system remains in its original energy state; as a result, the quantum state \( |0, +\rangle \) evolves adiabatically into the state \( |1, -\rangle \). Due to large value of the coupling strength close to the escape point, Landau-Zener transitions can be neglected. The final state, i.e. \( |1, -\rangle \) or \( |0, -\rangle \), is determined by the switching measurement.

VII. TUNABLE COUPLING

In this section, we calculate numerical values of the coupling strength \( g \), using the circuit parameters of Ref. 1 and show that the coupling can be tuned over a wide range with the bias parameters.

Fig. 16 presents the dependence of the coupling \( g \) as a function of the bias current \( I_b \) and the flux \( \Phi_S \), for the wells family \( [0] \) of the dc-SQUID. The solid line indicates where the two qubits are in resonance \( (\nu_S = \nu_T) \).
and a Cooper pair box \cite{32}. The coupling has been calculated using the expression (35) at the resonance \( \nu_S = \nu_T \) with the parameters of Ref. \cite{1} and only the ACPT asymmetries have been varied. For the asymmetric transistor with \( \lambda = \mu = -41.6\% \), the coupling \( g \) (red curve) is maximum at \( \delta = \pm \pi \) where it equals to 1217 MHz; it becomes zero at \( \delta = 0 \). In the case of the symmetric transistor with \( \lambda = \mu = 0 \), the coupling strength reads \( g = \alpha_\parallel / 2 \sqrt{E_C^{\parallel} / \nu_S} (\nu_S - \nu_T) \). Consequently, at the resonance, the coupling between a symmetric transistor and a SQUID is zero. For a Cooper pair box with \( E_{T2}^{\parallel} = 0 \) (\( \mu = -1 \) and \( \lambda = -41.6\% \)), the coupling reads \( g = \frac{\alpha_\parallel}{2} \sqrt{E_C^{\parallel} h \nu_S (1 + \lambda)} \). This corresponds to the result of Ref. \cite{31}. The calculated coupling \( g \) (black curve) does not depend on \( \delta \) and remains equal to 514 MHz.

Fig. 18(a) shows the dependence of the analytical and numerical couplings at the resonance as a function of the absolute value of the phase \( \delta \). The numerical simulations allow us to check the validity of the analytical coupling. These simulations have been realized by diagonalizing the full Hamiltonian in the basis of 20 charge states and the first 9 excited states of the dc-SQUID in the absence of anharmonicity. The dc-SQUID and ACPT Hamiltonians, as well as the Josephson and capacitive coupling Hamiltonians can all be expressed naturally in these bases. The numerical coupling is found by calculating the energy spectrum as a function of \( \delta \) between \(-\pi\) and \( \pi \) for a fixed frequency \( \nu_S \). The spectrum is first calculated without the coupling terms in order to find the two opposite values of \( \delta = \pm \delta_r \) where the ACPT and the dc-SQUID are in resonance (\( \nu_S = \nu_T \)), i.e. where the second and third energy bands intersect. In the presence of coupling, the degeneracy between the eigenstates is lifted and an anti-level crossing appears with an energy separation equal to the coupling \( g \) at \( \delta = \pm \delta_r \). As an example, the inset of Fig. 18(a) shows the energy spectrum for \( \delta_r = 0.5 \pi \) with and without coupling from which we extract \( g/\hbar = 251 \) MHz. The numerical and analytical simulations remarkably give quite close results which confirm the validity of the analytical expression (35). The theoretical coupling is calculated here without any free parameters by assuming equal Josephson and capacitive asymmetries \( \lambda = \mu = -41.6\% \). Finally, we note the good agreement with the experimental coupling measured in Ref. \cite{1} and shown in black points in Fig. 18. By adjusting the capacitive asymmetry to \( \lambda = -36.4\% \) the numerical coupling is in perfect agreement with the experiment as shown in Fig. 18(b).

VIII. CONCLUSION

In conclusion, we have analyzed in detail the theoretical quantum dynamics of a superconducting circuit based on a dc-SQUID in parallel to an Asymmetric Cooper Pair Transistor (ACPT). The Lagrangian of the circuit was first established from the current conservation equations expressed at each node of the circuit. The Hamiltonian, deduced from the Lagrangian, is decomposed in three distinct terms, namely the dc-SQUID, the ACPT and the coupling Hamiltonians. We first studied the individual dynamics of the dc-SQUID and the ACPT. Depending on its anharmonicity, the dc-SQUID can be seen either as a harmonic oscillator or as a phase qubit, whereas the ACPT behaves as a charge qubit. In addition to the optimal bias points \( (\delta = 0, n_g = 1/2) \) which was successfully demonstrated in a symmetric Cooper pair transistor, the ACPT presents a second optimal point \( (\delta = \pi, n_g = 1/2) \).
At these points, the charge qubit is insensitive in first order to the charge, flux and current noise and therefore shows a larger coherence time. We found that the coupling Hamiltonian between the dc-SQUID and the ACPT is made of two different terms corresponding to the Josephson and the capacitive couplings, which mix phases and charges of both sub-circuits, respectively.

The Hamiltonian of the full circuit was discussed through two different limits of the dc-SQUID. When it behaves as a harmonic oscillator, the quantum dynamics of the circuit is described by the well-known Jaynes-Cummings Hamiltonian. The microscopic circuit is then similar to a two-level atom coupled to a single-mode optical cavity. It offers compared to this latter a better tunability, a faster control and read-out of the quantum system, and a good scalability for complex architectures implementation. For example a circuit of several ACPT qubit in parallel could be considered whose quantum information will be mediated by the dc-SQUID [39]. When the anharmonicity of the dc-SQUID is strong, it behaves as a phase qubit. The full circuit is then described as the coupling of two different class of qubits, i.e., a phase and a charge qubit. The coupling Hamiltonian contains terms in \( \hat{\sigma}^T_y \hat{\sigma}^S_y \) and \( \hat{\sigma}^T_x \hat{\sigma}^S_x \) which are prominent when the two qubits are in resonance. These terms allow two-qubits gate operations as the \( \sqrt{\text{SWAP}} \) gate. In addition they enable to read out the quantum state of the ACPT by a nanosecond flux pulse as observed in Ref.[1]. Indeed such a pulse produces an adiabatic quantum transfer of the state \([0, +] \) into the state \([1, -] \), i.e. the energy quantum stored in the ACPT is transferred into the dc-SQUID in order to be detected. A non-resonant term in \( (\hat{a} - \hat{a}^\dagger) \hat{\sigma}^T_z \) or in \( \hat{\sigma}^S_y \hat{\sigma}^T_z \) is present in the Josephson coupling. Although its effect on the quantum dynamics is weak, this latter term explains the charge qubit read-out in the limit \( \nu_T \gg \nu_S \) by mean of an effective additional current in the dc-SQUID. That read-out method is employed in the Quantronium circuit [16].

In both limits of the dc-SQUID, we demonstrated that the coupling can be strongly tuned, mainly with the Josephson coupling term which has a strong phase \( \delta \) dependence. It can be used to accomplish two-qubit gate operations, and can also be turned off in order to perform one-qubit gate operations without disturbing the unaddressed qubit.

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Appendix A: Current conservation laws

The current conservation law, applied to each node of the circuit (see Fig. 1(b)), yields six equations for the active phases $\phi_1, \phi_2, \psi, \theta$ and $\xi$. These equations are identical to the Euler-Lagrange equations derived from the circuit Lagrangian, (2), (3), (4) and read, respectively,

\begin{align*}
\phi_0 \frac{\dot{\xi} - \phi_1}{L_1} - C_0 \phi_0 \dot{\phi}_1 - I_0 \sin(\phi_1) &= 0; \\
\phi_0 \frac{\dot{\gamma} - \phi_2}{L_2} - C_0 \phi_0 \dot{\phi}_2 - I_0 \sin(\phi_2) &= 0; \\
I^2_2 \sin(\xi - \psi + \phi_T) + C_2^T \phi_0 (\xi - \phi_T) - C_g \phi_0 (\dot{\psi} - \dot{\phi}_T) - I^2_2 \sin(\xi - \psi + \phi_T) &= 0; \\
C_g \psi - (C_P + C_g) \dot{\theta} &= 0; \\
\phi_0 \frac{\dot{\gamma} - \phi_S - \phi_0}{L_3} - C_2^T \phi_0 (\dot{\xi} - \phi_T) - I^2_2 \sin(\xi - \psi + \phi_T) &= 0.
\end{align*}

In Sec. II C we show that this system of six equations can be reduced to four equations by ignoring the high frequency quantum dynamics of the phases $\gamma$ and $\xi$.

Appendix B: Parameters of the coupled circuit studied in Ref. 1

Throughout this article, we illustrate the theory with numerical values and plots calculated by using the parameters of the circuit studied in Ref. 1. These parameters are collected in Tab. III.

Appendix C: Conjugate variables

The phases $(X_{||}, Y_{||}, \psi, \theta)$ and their conjugate momenta $(-hP_{||}, -hP_{\perp}, -hn, -hn_Q)$ are the appropriate variables of the circuit Hamiltonian 15. The momenta are related to the velocities $(\dot{X}_{||}, \dot{Y}_{||}, \dot{\psi}, \dot{\theta})$ involved in the kinetic part of the Lagrangian 13 by the following expressions 22.

| Parameters of the dc-SQUID | Label | Value |
|-----------------------------|-------|-------|
| Critical current per Josephson junction (JJ) | $I_0$ | 1.356 $\mu$A |
| Capacitance per JJ | $C_0$ | 227 fF |
| Loop inductance | $L_s$ | 190 pHz |
| Inductance asymmetry | $\eta$ | 0.28 |
| Bidimensionality parameter | $b$ | 1.28 |

TABLE III. Electric parameters of the coupled circuit studied by Fay et al. 1.

\begin{align*}
-h P_{||} &= \frac{\partial L}{\partial X_{||}} = \phi_0^2 \left\{ (2C_0 + \alpha || C_2^T) \dot{X}_{||} - \alpha || C_2^T \left( \alpha \dot{Y}_{||} + \dot{\psi} \right) \right\} \\
-h P_{\perp} &= \frac{\partial L}{\partial Y_{||}} = \phi_0^2 \left\{ (2C_0 + \alpha \perp C_2^T) \dot{Y}_{||} - \alpha \perp C_2^T \left( \alpha || \dot{X}_{||} - \dot{\psi} \right) \right\} \\
-h n &= \frac{\partial L}{\partial \psi} = \phi_0^2 \left\{ C_2 \dot{\psi} - C_2^T (\alpha || \dot{X}_{||} - \alpha \perp \dot{Y}_{||}) - C_g \dot{\theta} \right\} \\
-h n_Q &= \frac{\partial L}{\partial \theta} = \phi_0^2 \left\{ (C_P + C_g) \dot{\theta} - C_g \dot{\psi} \right\}.
\end{align*}

Appendix D: Pauli matrices

The Pauli matrices related to the dc-SQUID are defined in the eigenbasis of the phase qubit $\{ |1\rangle, |0\rangle \}$ as

\begin{align*}
\sigma_z &= \begin{pmatrix} 1 & 0 \\
0 & -1 \end{pmatrix}, \quad \sigma_x = \begin{pmatrix} 0 & 1 \\
1 & 0 \end{pmatrix}, \quad \sigma_y = \begin{pmatrix} 0 & -i \\
i & 0 \end{pmatrix}.
\end{align*}

(D1)

Similarly the Pauli matrices related to the ACPT (\(\sigma_z^T\), \(\sigma_x^T\), \(\sigma_y^T\)) are defined in the eigenbasis of the charge qubit $\{ |+\rangle, |-\rangle \}$.
