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Quantum computing with exciton-polariton condensates

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Exciton-polariton condensates have attractive features for quantum computation, e.g., room temperature operation, high dynamical speed, ease of probe, and existing fabrication techniques. Here, we present a complete theoretical scheme of quantum computing with exciton-polariton condensates formed in semiconductor micropillars. Quantum fluctuations on top of the condensates are shown to realize qubits, which are externally controllable by applied laser pulses. Quantum tunneling and nonlinear interactions between the condensates allow SWAP, square-root-SWAP and controlled-NOT gate operations between the qubits.

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INTRODUCTION

Exciton-polaritons are hybrid light-matter particles, where the small dephasing of a photonic system is blended with the particle–particle interactions from a condensed matter system. The overwhelming majority of research into exciton-polaritons has been focused on semi-classical physics, that is, physics where the quantization of the particle number according to quantum theory is not itself apparent: Bose-Einstein condensation and polariton lasing;[1] quantum fluids and solitons;[2] and topological polaritons[3] are all well-described within semi-classical theory. Nevertheless, it has always been expected that exciton-polaritons are ultimately quantum particles,[4] theoretically capable of entanglement,[5,6] and it has been speculated that polaritons could perform as quantum computers[7,8] or quantum simulators.[9,10]

As the quality factor of microcavities has steadily progressed, there has been a recent resurgence of interest in quantum exciton-polariton physics. Through exciting a polariton with an entangled photon it was proven that exciton-polaritons are indeed quantum particles, capable of carrying correlations at the quantum level.[11] The onset of the polariton blockade[12] has also shown that the repulsive interaction energy between just two polaritons is considerable,[13,14] and that suggestions of strongly interacting polaritons can be reachable in the future.[15,16]

However, serious challenges must be addressed before claiming that exciton-polaritons can perform as a quantum computer. First, a suitable basis for encoding information must be chosen and it must be shown that exciton-polaritons are compatible with such a basis. Typically, architectures for quantum computation are based on qubits, which require a two level quantum system. It is not immediately obvious that exciton-polaritons can serve in this regard, since they typically form multi-particle states. Even though binary degrees of freedom (e.g., spin or vorticity[17,18]) exist, this does not automatically mean that polaritons can operate as qubits as the different degrees of freedom can themselves be populated by a range of particle number (Fock) states.

Even if a quantum computational basis can be established, a significant task remains in designing quantum logic gates (assuming that one aims for emulation of the most well-known quantum circuit model of quantum computation). If a set of universal quantum gates can be established, then in principle any unitary operation can be reproduced on a given input state. For example, a set of single qubit logic gates combined with square-root-SWAP gate (or controlled-NOT gate) is universal. However, it is essential that the gates can be arbitrarily connected and formed into scalable circuits, which is already a challenging task for classical polaritonic logic gates.[19]

Finally, it is important that the logic gates should be operable on a timescale faster than the dissipation and dephasing rates of exciton-polaritons. While classical information carried in the form of a polariton wavepacket can be re-amplified with a non-resonant external laser,[20] which replaces lost particles with those of the same classical properties (e.g., phase and polarization), applying the same to a quantum polariton state actually causes quantum information to be lost faster.[21] Essentially, if a polariton carrying quantum correlations is lost from the system it can not be replaced with a fully identical particle. Attempting to do so necessarily worsens the situation as the replacement particle can not possibly have the same quantum correlations with the rest of the system as the particle that was lost. Thus all gates need to act before polaritons escape the system. Fortunately, state-of-the-art microcavities support polariton lifetimes exceeding hundreds of picoseconds,[22] while applied laser pulses can be operated on sub picosecond scale or faster.

Here we address the above challenges from the theoretical point of view, considering exciton-polariton condensates in micropillar arrays[9] as a platform. Even though a polariton condensate is composed of many particles, we show that number fluctuations on top of the mean-field value themselves correspond to an anharmonic oscillator in the presence of moderate polariton–polariton interactions. Transitions between the lowest two energy levels are non-resonant with other levels, effectively allowing a qubit basis, which can be measured with homodyne detection techniques. As outlined in Fig. 1, single qubit gates are controllable by laser parameters and the availability of quantum tunneling between neighboring micropillars, it is possible to realize SWAP and square-root-SWAP (sSWAP) gates, which allow for moving qubits between different pillars and universal quantum operations to establish a complete scalable quantum circuit architecture. Additionally, the availability of cross–Kerr interactions between polaritons, which are allowed via the presence of cross–spin interactions, are shown to realize a quantum controlled-NOT (cNOT) gate.

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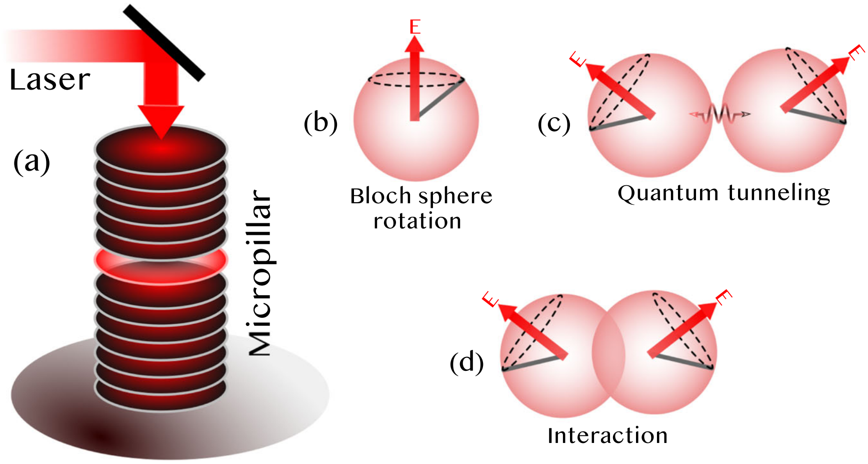

Fig. 1 A scheme of quantum computing with exciton-polariton condensates. a A polariton condensate confined in a semiconductor micropillar and excited with a coherent laser supports a qubit suitable for quantum computing. **b** The evolution under the effective qubit Hamiltonian induces rotation on the Bloch sphere. The amount of rotation and the axis of rotation is controlled by the laser parameters. With specific sets of laser parameters we achieve single-qubit quantum gates. **c** SWAP and sSWAP gates: quantum tunneling between two micropillars allows SWAP and sSWAP (square-root-SWAP) operations between qubits in different micropillars. **d** cNOT gate: accounting for the spin degree of freedom, two qubits can be encoded in the single micropillar and interact via a cross-Kerr type interaction to induce the cNOT gate operation.

THE EXCITON-POLARITON CONDENSATE

We consider a semiconducting micropillar that supports a nonlinear polariton mode $\hat{a}$. The mode is excited with a coherent optical field $\mathcal{P}(t)$. The Hamiltonian is thus

$$H = \Delta\ \hat{a}^{\dagger}\hat{a} + \alpha\ \hat{a}^{\dagger}\hat{a}^{\dagger}\hat{a}\hat{a} + \mathcal{P}(t)^{*}\hat{a}^{\dagger} + \mathcal{P}(t)\hat{a}, \tag{1}$$

where Δ is a detuning between the optical field and the polariton mode and α is the Kerr nonlinear interaction strength. We define a condensate as a coherent state, which is readily generated by the application of a constant coherent optical field $\mathcal{P}(t) = P_0$ with the occupation number, $N_c = \langle\hat{a}^{\dagger}\hat{a}\rangle$. We aim to engineer the quantum operations by switching the optical pump from $\mathcal{P}(t) = P_0$ to $\mathcal{P}(t) = Pe^{i\varphi}$ for a very short time compared to the polariton lifetime. Since the condensate formed by P_0 does not have enough time to decay, the short excitation $Pe^{i\varphi}$ will only induce fluctuations on top of the preformed condensate. We now consider the following change of variable: $\hat{a}^{\dagger} = \sqrt{\hat{N}}\exp(i\hat{\theta})$, where $\hat{N}$ and $\hat{\theta}$ are the number and phase operators of the condensate satisfying the commutation relation $[\hat{\theta}, \hat{N}] = i$. In terms of these operators the Hamiltonian reads:

$$H = \Delta\hat{N} + \alpha\hat{N}(\hat{N} - 1) + 2P\sqrt{\hat{N}}\cos(\varphi - \hat{\theta}) \tag{2}$$

The polariton number operator can be expressed as $\hat{N} = \hat{a}^{\dagger}\hat{a} = N_c + \hat{n}$, where $N_c = \langle\hat{a}^{\dagger}\hat{a}\rangle$ is the mean-field part of the condensate and $\hat{n}$ is the quantum part. Note that the mean field N_c cannot carry quantum information itself. Instead, the quantum part $\hat{n}$ will be used for encoding quantum information. We express the Hamiltonian in terms of $\hat{n}$:

$$H = C + \Omega\ \hat{n} + \alpha\hat{n}^2 + 2P\sqrt{(N_c + \hat{n})}\quad \cos(\varphi - \hat{\theta}), \tag{3}$$

where $C = \Delta N_c + \alpha N_c(N_c - 1)$ is a classical number and $\Omega = \Delta + \alpha(2N_c - 1)$ is defined as the effective detuning. Here we now consider that the magnitude of N_c is much larger than the quantum fluctuations $\hat{n}$ (which is the standard case for polariton condensates), such that the effective Hamiltonian is given by,

$$H_f = \Omega\hat{n} + \alpha\hat{n}^2 + 2P\sqrt{N_c}\quad \cos(\varphi - \hat{\theta}) \tag{4}$$

Note that inside the square-root, we have neglected the operator $\hat{n}$ which is justified when $\sqrt{\langle\hat{n}^2\rangle / N_c} \ll 1$. In Fock space, the number operator satisfies $\hat{n}|n\rangle = n|n\rangle$ and the corresponding phase operator satisfies $\exp(\pm i\hat{\theta})|n\rangle = |n \pm 1\rangle$ (see the Supplementary Information). The effective Hamiltonian in the Fock space then reads,

$$H_f = \sum_n \left[(\Omega n + \alpha n^2)|n\rangle\langle n| + P\sqrt{N_c}\left(e^{i\varphi}|n\rangle\langle n+1| + e^{-i\varphi}|n+1\rangle\langle n|\right)\right] \tag{5}$$

This effective system is formed with an exciton-polariton condensate in a micropillar interacting with a laser field. Note that a coherent condensate can be achieved even in the strongly interacting regime (see Supplementary Information). The system has similarity with a superconducting qubit formed in a Josephson junction between two superconducting islands. While one among the two superconducting islands is analogous to the condensate in the micropillar, the second one is replaced by the classical laser field which can be manipulated (through controlling its phase, detuning and amplitude) with much more ease and accuracy than a cryogenically cooled superconducting island (controlling phase, energy and amplitude of the superconducting wavefunction are much more challenging). Moreover, the condensate in a semiconductor micropillar itself can operate at much higher temperatures (even at room temperature) than the operating temperatures of superconducting qubits.

ANHARMONIC SPECTRUM AND THE QUBIT

In Fig. 2, we find that the spectrum of the Hamiltonian is periodic in the effective detuning Ω with a period 2α. Let us replace Ω with an effective detuning parameter ω that varies between $\pm\alpha$ such that the spectrum is restricted within one period. Moreover, in Fig. 2 we see that the gaps between the lowest two energy levels and the next ones are different. According to our need, we can operate in some regime where the difference in the gaps can be obtained by diagonalizing the Hamiltonian given by Eq. (5). These unequal gaps are originating from the anharmonic behavior of the Hamiltonian, which can be explicitly seen in Eq. (4) in the form of the cosine term (if instead of the cosine term one had a $\hat{\theta}^2$ term, one would have a harmonic oscillator). Let us consider that δE_1 is the gap between the first two energy levels and δE_2 is the same between the second and third energy levels. In the regime, where $|\delta E_1 - \delta E_2|$ is larger than the linewidth γ, we can consider only the lowest two energy levels as our qubit basis $|0\rangle$ and $|1\rangle$. In this low

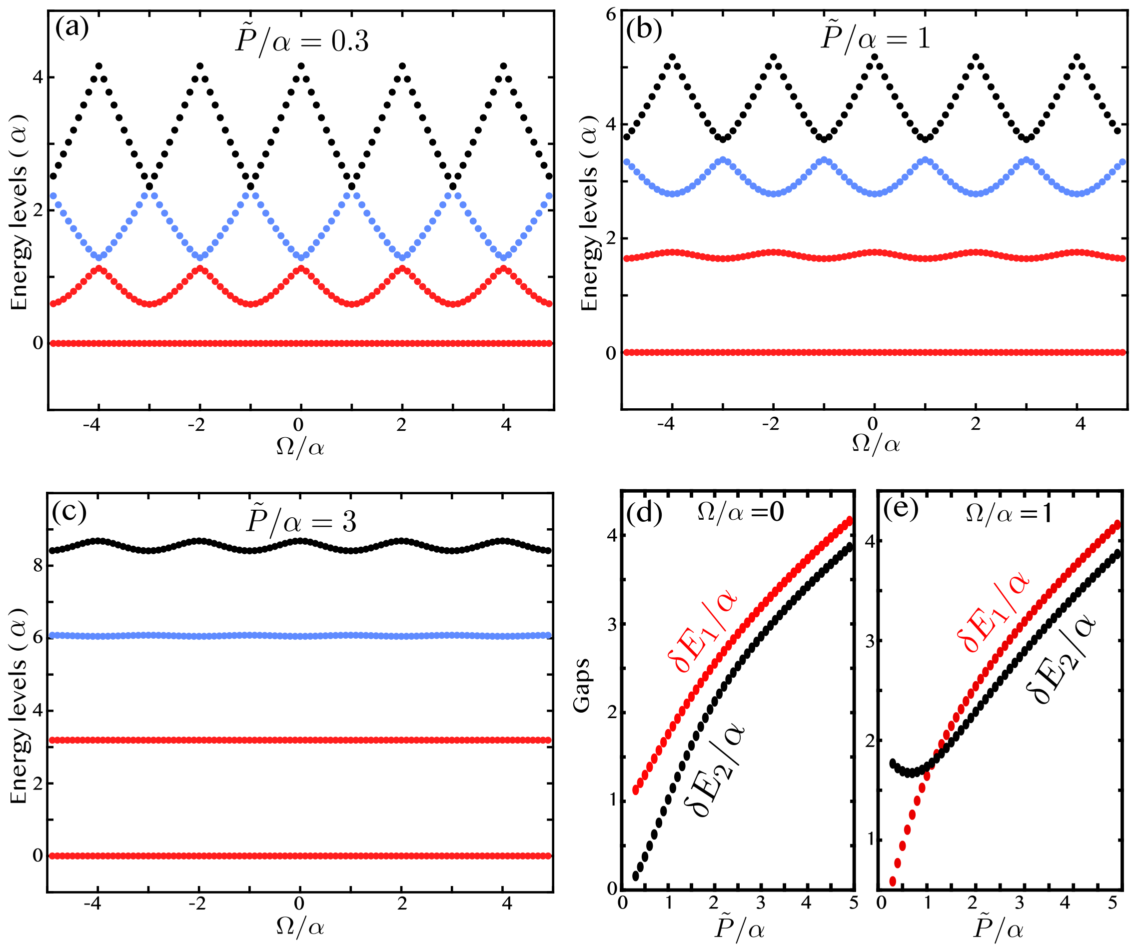

Fig. 2 Spectrum of the Hamiltonian in different regimes. We show the lowest four energy levels as functions of Ω (the effective detuning) for $\tilde{P}/\alpha<1$ (**a**), $\tilde{P}/\alpha=1$ (**b**), and $\tilde{P}/\alpha>1$ (**c**), where $\tilde{P}=P\sqrt{N_c}$. In all regimes, the spectrum is periodic in Ω with a period 2α. In **d** and **e**, we show the two energy gaps δE_1 and δE_2 between lowest energy levels as functions of $\tilde{P}/\alpha$ at two different values of the effective detuning Ω. Here δE_1 is the gap between the first two levels and δE_2 is the same between the second and third levels.

energy qubit subspace, the effective qubit Hamiltonian reads $\hat{H}_q=E_x\hat{\sigma}_x+E_y\hat{\sigma}_y+E_z\hat{\sigma}_z$, where $\hat{\sigma}_i$ are the Pauli matrices, $E_x=P\sqrt{N_c}\cos\varphi$, $E_y=P\sqrt{N_c}\sin\varphi$, and $E_z=(\alpha-\omega)/2$. Note that here we have excluded the overall energy shift $\mathbb{1}(\alpha-\omega)/2$ from the Hamiltonian. Formally, this Hamiltonian can be rewritten in a compact form:

$$\hat{H}_q=\boldsymbol{E}\cdot\hat{\boldsymbol{\sigma}}, \tag{6}$$

where the vectors $\mathbf{E}=(E_x, E_y, E_z)$ and $\hat{\boldsymbol{\sigma}}=(\hat{\sigma}_x,\hat{\sigma}_y,\hat{\sigma}_z)$. This qubit Hamiltonian can be externally controlled by tuning the vector **E** through controlling the laser detuning ω, strength P, and the phase φ.

SINGLE-QUBIT GATES

For universal quantum computing, single-qubit gates are required. A single-qubit gate operation is equivalent to a unitary rotation of a qubit around a given axis. Here, we show that the evolution of an initially prepared state under the effective Hamiltonian $\hat{H}_q$ can induce an arbitrary single-qubit quantum gate operation. For specific operation, the parameters involved in the Hamiltonian must be chosen appropriately. If the duration of the application of the pump $Pe^{i\varphi}$ is τ, then the evolution operator is given by $\exp(-i\tau\hat{H}_q/\hbar)$, which can be rewritten as,

$$U_\epsilon(\beta)=\mathbb{1}\ \cos\beta-i\boldsymbol{\epsilon}.\hat{\boldsymbol{\sigma}}\sin\beta, \tag{7}$$

where $\beta=\tau E/\hbar$, and $\boldsymbol{\epsilon}=\boldsymbol{E}/E$ is a unit vector controlled by the system parameters, and $E=\sqrt{P^2N_c+(\alpha-\omega)^2/4}$. This is a general form of a unitary matrix in single-qubit space. More explicitly the operator can be written in matrix form:

$$U_\epsilon(\beta)=\begin{bmatrix}\cos\beta-i\sin\beta\cos\xi & -i\sin\beta\sin\xi\ e^{-i\varphi}\\ -i\sin\beta\sin\xi\ e^{i\varphi} & \cos\beta+i\sin\beta\cos\xi\end{bmatrix}, \tag{8}$$

where we parameterized the unit vector in angular coordinates $\boldsymbol{\epsilon}=(\sin\xi\cos\varphi,\ \sin\xi\sin\varphi,\ \cos\xi)$. By choosing appropriate β, ξ, and φ we can obtain different quantum operations, e.g., a rotation $\hat{R}_x(\beta)$ around $\hat{\sigma}_x$ with ($\xi=\pi/2$, $\varphi=0$), and a rotation $\hat{R}_y(\beta)$ around $\hat{\sigma}_y$ with ($\xi=\pi/2$, $\varphi=\pi/2$). Since any single qubit unitary operation can be composed as $\hat{R}_x(\beta_1)\hat{R}_y(\beta_2)\hat{R}_x(\beta_3)$ (where $\beta_{1,2,3}$ correspond to three time durations),[23] arbitrary single-qubit gates can be obtained with a fixed $\xi=\pi/2$ and controlling only β and φ. This shows that arbitrary single-qubit operations are allowed by a pulse with a fixed detuning $\omega=\alpha$ (equivalently $E_z=0$) but a time dependent phase φ. However, additional control on the laser field amplitude P provides control on ξ, which allows more flexibility on achieving single-qubit gates, e.g., a frequently used single-qubit gate, known as the Hadamard gate, can be readily realized with ($\beta=\pi/2$, $\xi=\pi/4$, $\varphi=0$).

SWAP AND SSWAP GATES

Let us now consider that two physically separated polariton condensates with associated field operators $\hat{a}_1$ and $\hat{a}_2$ are coupled via a coherent (Josephson) tunneling term. The corresponding Hamiltonian is given by,

$$H_{12} = \sum_{j=1,2} H_j + J(\hat{a}_1^\dagger \hat{a}_2 + \hat{a}_2^\dagger \hat{a}_1), \tag{9}$$

where $H_j = \Delta_j \hat{a}_j^\dagger \hat{a}_j + \alpha_j \hat{a}_j^\dagger \hat{a}_j^\dagger \hat{a}_j \hat{a}_j + P_j(\hat{a}_j e^{i\varphi_j} + \hat{a}_j^\dagger e^{-i\varphi_j})$ corresponds to a single condensate. Following the same procedure as we have shown for a single condensate, we obtain a low energy effective Hamiltonian:

$$H_{q_1 q_2} = \sum_{j=1,2} \boldsymbol{E}^j \cdot \hat{\boldsymbol{\sigma}}^j + E_T(\hat{\sigma}_+^1 \hat{\sigma}_-^2 + \hat{\sigma}_+^2 \hat{\sigma}_-^1), \tag{10}$$

where E_T is the tunneling amplitude. This Hamiltonian allows both efficient SWAP and sSWAP operations. While a SWAP operation swaps the quantum states between the two quibts, sSWAP operations establish universality (any quantum operation can be achieved with sSWAP and single-qubit gates).

Our considered system is a driven dissipative system. Exciton-polaritons have a finite lifetime $\hbar/\gamma$. An accurate description of this driven dissipative system is given by the quantum master equation:

$$i\hbar \dot{\hat{\rho}} = [\hat{H}_{q_1 q_2}, \hat{\rho}] + \frac{\gamma}{2} \sum_{j=1,2} \mathcal{L}^j(\hat{\rho}), \tag{11}$$

where $\hat{\rho}$ is the density matrix describing the qubits, $\mathcal{L}^j(\hat{\rho}) = 2\hat{\sigma}_-^j \hat{\rho} \hat{\sigma}_+^j - \hat{\sigma}_+^j \hat{\sigma}_-^j \hat{\rho} - \hat{\rho} \hat{\sigma}_+^j \hat{\sigma}_-^j$, and $\hat{\sigma}_\pm^j$ are the raising and lowering operators. In Fig. 3b, we show the gate fidelities of sSWAP and SWAP gates as the function of the ratio of the polariton–polariton interaction strength to the dissipation rate.

For efficient gate operations we need to externally control the tunneling amplitude E_T, such that the quantum tunneling can be activated at will. In principle, Josephson coupling can be modulated by applying stress[24] or a magnetic field,[25] however, here we opt for a scheme where two qubits in two physical semiconductor micropillars are connected by a third micropillar placed between them (see the inset of Fig. 3b). This three site system is considered with a fixed nearest-neighbor tunneling amplitude, g (see Fig. 4). The on-site energy of the middle site could be controlled via an external laser field. In Fig. 4, we show the probability of the transition $|10\rangle \to |01\rangle$ as a function of the ratio between the relative detuning Δ_s and the tunneling

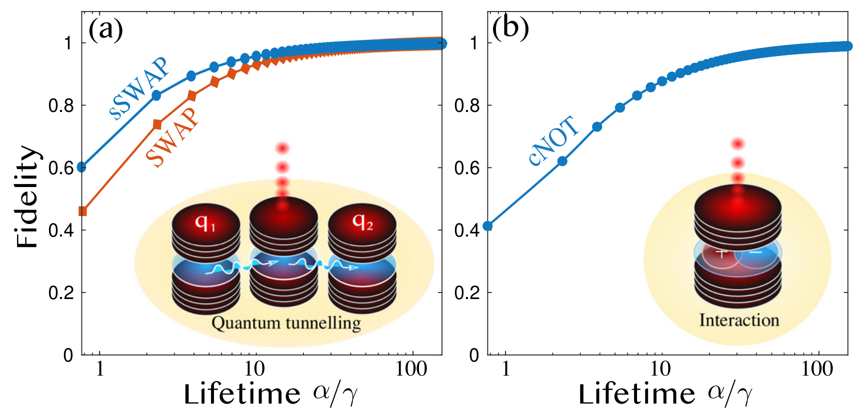

Fig. 3 Two-qubit gate fidelities as functions of the ratio between the polariton-polariton interaction strength and the dissipation rate (α/γ). a Fidelities of square-root-SWAP (blue circles) and SWAP (red diamonds) gates as functions of polariton lifetime. The inset shows a scheme for sSWAP and SWAP operations based on three coupled micropillars. For sSWAP gate, we considered $E_x^j/E_z^j = 11.55$, $E_y^j/E_z^j = 0$, $E_z^j/\alpha = 1$, $E_T/\alpha = 45.64$ and pulse duration $\tau = 0.576\hbar/\alpha$ and for SWAP gate, we considered $E_x^j/E_z^j = 7.6$, $E_y^j/E_z^j = 0$, $E_z^j/\alpha = 1$, $E_T/\alpha = 48.7$ and pulse duration $\tau = 1.08\hbar/\alpha$. **b** Fidelity of a maximally entangled state obtained with a cNOT gate as a function of the polariton lifetime. We considered the parameters $E_x^1 = E_z^1 = 0$, $E_x^2/E_z^2 = 0.05$, $E_z^2/\alpha = 1$, $E_z^{12}/\alpha = \alpha_{12}/\alpha = 1$ and pulse duration $\tau = 1.36\hbar/\alpha$ (here $\alpha_j = \alpha$ is the nonlinearity strength). The inset shows a scheme for cNOT operation based on two spin components of a polariton condensate.

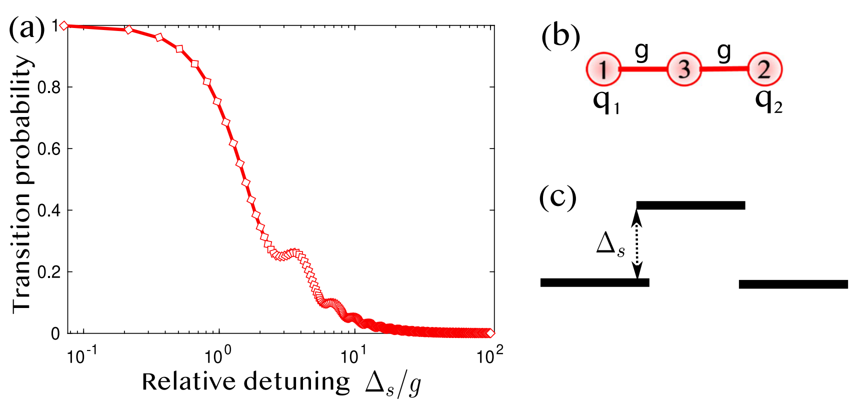

Fig. 4 Controlling the tunneling amplitude by tuning the detuning. a Probability of the transition $|10\rangle \to |01\rangle$ between the qubits as a function of the ratio between the relative detuning Δ_s and the tunneling amplitude g between the nearest-neighbor sites. Here the transition occurred with a unitary evolution for a time duration τ. **b** A three site lattice system with nearest-neighbor tunneling amplitude g. **c** The site energies: the site at the middle has a relative energy Δ_s with respect to the qubit sites placed at the two ends. We consider $g/\alpha = 2.07$, and a time duration $\tau = 1.08\hbar/\alpha$.

amplitude g between the neighboring micropillars. We find that the transition probability can be tuned from a high value to a low one by tuning the parameter Δ_s/g.

CONTROLLED-NOT GATE

Although a cNOT operation can be obtained as a combination of single-qubit and sSWAP gates, here we show that a direct realization of cNOT gate is also possible in semiconductor micropillars. Let us consider that two polariton condensates with associated field operators $\hat{a}_1$ and $\hat{a}_2$ are interacting with a cross-Kerr type nonlinearity. The corresponding Hamiltonian is given by,

$$H_{12} = \sum_{j=1,2} H_j - 2\alpha_{12}\ \hat{a}_1^\dagger\hat{a}_2^\dagger\hat{a}_2\hat{a}_1, \tag{12}$$

where the Hamiltonian $H_j = \Delta_j\hat{a}_j^\dagger\hat{a}_j + \alpha_j\hat{a}_j^\dagger\hat{a}_j^\dagger\hat{a}_j\hat{a}_j + P_j(\hat{a}_j e^{i\varphi_j} + \hat{a}_j^\dagger e^{-i\varphi_j})$ describes a single condensate, α_j and α_{12} are the Kerr and cross-Kerr nonlinearity strengths, and P_j, φ_j, and Δ_j are the amplitudes, phases and detuning values of the applied laser fields. The cross-Kerr nonlinearity can be physically realized with exciton-polaritons by considering the spin degree of freedom. Exciton-polaritons can have one of two spin projections on the structure growth axis,[26] corresponding to two independent quantum modes. Noise in the spin is drastically suppressed in the regime of polariton condensates.[27] While the quantum gates in our schemes are aimed to operate in a fraction of the polariton lifetime, timescales associated to spin noise and the number fluctuations are much larger for large condensate occupation.[27] It has also been well established that cross-Kerr type nonlinearity is present between the opposite spin states, and a variety of interaction strengths have been achieved experimentally.[28] Alternatively, cross-phase modulation between neighboring optical cavity modes has been discussed in different works,[29,30] although not yet realized experimentally for exciton-polariton systems.

Following the same procedure as we have shown for a single condensate, we obtain a low energy effective Hamiltonian:

$$H_{q_1q_2} = \sum_{j=1,2} \boldsymbol{E}^j.\hat{\boldsymbol{\sigma}}^j - E_z^{12}\ (\mathbb{1} + \sigma_z^1)\ (\mathbb{1} + \sigma_z^2), \tag{13}$$

where $\boldsymbol{E}_j = [P_j\sqrt{N_c}\cos\varphi_j,\ \ P_j\sqrt{N_c}\sin\varphi_j,\ \ (\alpha_j - \omega_j)/2]$, $E_z^{12} = \alpha_{12}/2$, and ω_j varies between $\pm\alpha_j$. For demonstrating cNOT operations, we consider the two-qubit Hamiltonian $H_{q_1q_2}$ with $\varphi^j = 0$. It has been shown earlier that this Hamiltonian can be used for inducing cNOT operation.[31] For instance, under a pulse of duration τ, the evolution operator is given by,

$$U_{q_1q_2} = e^{-i\frac{\tau}{\hbar}H_{q_1q_2}}. \tag{14}$$

We find that this evolution operator becomes a cNOT operation for $E_z^1 = 0$, $E_y^2/E_z^2 \ll 1$, and a suitable τ. In Fig. 3b, we show the fidelity of an exciton-polariton cNOT gate as a function of the ratio of the polariton-polariton interaction strength to the dissipation rate.

DISCUSSION

We have arrived at a theoretical scheme for implementation of the quantum circuit model of quantum computation using exciton-polariton condensates. Qubits are encoded in an anharmonic oscillator formed by the quantized fluctuations of the particle number about the mean-field value. A universal set of quantum gates (single-qubit and sSWAP gates) can be engineered using individual micropillars and quantum tunneling among them. SWAP gates are introduced to allow the coherent transport of quantum information between qubits. We further showed that cNOT gates can be directly induced by exploiting the available cross-Kerr interactions between modes with different polarizations. Other quantum gates, e.g., iSWAP and square-root-iSWAP can also be implemented straightforwardly within our scheme (see Supplementary Information). These gates would be controllable via the use of external laser pulses.

Our scheme is best suited for condensates with large number of polaritons. While a condensate with 50–100 polaritons can induce high quality quantum gates, smaller polariton occupancies introduce losses in the gate fidelity (see Supplementary Information). Technological imperfections can also introduce effects like small polarization splitting between the spin modes in a micropillar. However, the gate fidelity remains almost unaffected for splittings smaller than the polariton–polariton interaction strength (see Supplementary Information).

Aside exciton–exciton interaction another source of nonlinearity comes from the saturation of the exciton–photon coupling, which can be mapped to a renormalized Kerr nonlinear term in our considered Hamiltonian in Eq. (1).[2] Similarly, the biexciton nonlinearity can also be accounted for by renormalizing the interaction constant between polaritons with opposite spins.[32]

While exciton-polariton condensates allow high fidelity quantum gates for lifetimes much larger than the gate operation time, limited lifetime can introduce errors in the quantum circuits. However, fault-tolerant quantum computers can be achieved by encoding a logical qubit in multiple exciton–polariton qubits (allowing quantum error correction).[23,33–35] In principle, dynamical decoupling schemes appropriate for open quantum systems could also extend the effective loss time.[36] Finally, it is notable that typical polariton micropillar lattices can be fabricated with micron scale precision,[3] while the coherence of polariton condensates has been reported extending over a fraction of a millimeter.[37] The state-of-the-art in quantum computing has been developing rapidly in recent years, with companies developing systems with around 50 qubits,[38] while superconducting qubit systems[39] and ion trap[40] systems have already achieved 8 and 20 qubits, respectively. We hope that polariton lattices, which can potentially have a size of around $100 \times 100 = 10^4$ qubits (or double accounting for spin), will also be seen as relevant candidates.

We note that typically a fraction of a milliwatt of laser power is needed to coherently excite a micron sized condensate. A typical $1\ \text{mm}^2$ sized sample, which could contain 10^6 individual condensates, would require almost kilowatt power to achieve cryogenic liquid Helium temperature. As superconducting qubits occupy a much larger area, their cryogenic power required per qubit is much higher. In addition, exciton-polaritons can operate faster and at higher temperature.

DATA AVAILABILITY

The numerical data presented in this study is available from the authors upon reasonable request.

CODE AVAILABILITY

The codes for solving the quantum master equation for coupled mode systems are available from the authors upon reasonable request.

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AUTHOR CONTRIBUTIONS

S.G. conceived the project through discussions with T.L. S.G. performed the theoretical calculations. Both authors interpreted the results and wrote the manuscript. T.L. supervised the project.

COMPETING INTERESTS

The authors declare no competing interests.

ADDITIONAL INFORMATION

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