Proposal for Microwave Boson Sampling

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Microwave Boson Sampling

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Boson Sampling, the task of sampling the probability distribution of photons at the output of a photonic network, is believed to be hard for any classical device. Unlike other models of quantum computation that require thousands of qubits to outperform classical computers, boson sampling requires only a handful of single photons. However, a scalable implementation of Boson Sampling is missing. Here we show how superconducting circuits provide such platform. Our proposal differs radically from traditional quantum-optical implementations: Rather than injecting photons in waveguides, making them pass through optical elements like phase shifters and beam splitters and finally detecting their output mode, we prepare the required multi-photon input state in a superconducting resonator array, control its dynamics via tunable and dispersive interactions and measure it with non-demolition techniques.

Introduction.- The first post-classical computation will most probably be performed not on a universal quantum computer, but rather on a dedicated quantum hardware. A strong candidate for achieving this is represented by the task of sampling from the output distribution of a linear quantum optical network. This problem, known as boson sampling, has recently been shown to be intractable for any classical computer [1]. Aaronson and Arkhipov demonstrated that sampling the distribution of photons at the output of a linear photonic network is computationally inefficient for any classical computer, since it would require the estimation of an insurmountable amount of matrix permanents [2]. Such hardness proof is remarkably important since it shows that intermediate quantum setups can challenge the extended Church-Turing (ECT) thesis by suggesting a physical implementation that computes more efficiently than a nondeterministic Turing machine. In practice, the ECT thesis is not directly refutable since it refers to an asymptotically large scale implementation of a physical device, but the clear indication of a scalable setup and the neat experimental demonstration of such computation in medium-size devices would constitute a serious indication to reconsider the ECT thesis.

The emphasis of the previous argument points to the scalability issue. In fact, the original boson sampling setup works with optical photons that are difficult to generate as single photons in a deterministic way and that, given the state-of-the-art, cannot be detected with almost unit efficiency . Subsequent proposals have suggested the use of different initial states, like two-mode squeezed states [3, 4], photon added/subtracted coherent states [5], or vacuum squeezed states [6]. These modifications only partially solve the bottlenecks of nondeterministic state preparation and detection efficiency making the actual implementation of boson sampling exponentially demanding in the number of photons [7, 8]. A different approach to overcome such difficulties is the use of alternative experimental setups, such as trapped ions [9]. Unfortunately, the required interactions are not the natural ones for the setup considered, so frequent and localized laser pulses are necessary to constantly alter the dynamics with active control techniques.

In this Letter, we propose the realization of boson sampling with photons in the microwave regime by using superconducting circuits. We show how microwave photons are ideal for a scalable implementation satisfying the following three fundamental requirements of the problem: I) deterministic state preparation, II) direct implementation of the appropriate many-boson dynamics, and III) highly efficient photon-number measurements. In our proposal, we substitute the open-end optical waveguides with identical superconducting resonators, one for each mode, and couple them through a superconducting ring coupler implementing a tunable beam splitter Hamiltonian. Phase shifters are naturally implemented by tuning the resonator frequency in an independent way with the aid of an adjacent superconducting qubit. In this setup, deterministic state preparation can be efficiently prepared by loading the corresponding state of the qubit into each resonator using the Jaynes-Cummings interaction in circuit quantum electrodynamics (circuit QED) [10, 11]. The introduction of additional low-quality-factor (low-Q) resonators allows the system readout through a quantum non-demolition measurement. All of the required operations can be performed deterministically and with high fidelity on state-of-the-art superconducting devices [12–15]. This guarantees the scalability of our proposal and suggest superconducting platforms as a major physical candidate to the realization of large scale boson sampling experiments.

Finally, the advantages of the proposed implementation do not only help one address computational complexity questions alone, even if of primary importance. Recently, a modified version of the original apparatus has been shown to be the essential component of quantum

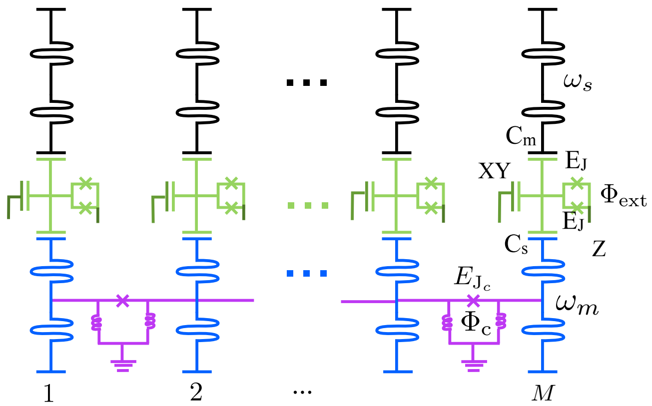

FIG. 1. Experimental proposal for scalable boson sampling with microwave photons. Cross-shaped transmon qubits (green) are capacitively coupled to a high-Q resonator (blue) and a low-Q measurement resonator (black). An XY control line allows single qubit rotations, while a Z control line changes the qubit frequency through the external flux $\Phi_{\rm ext}$. A superconducting ring intersected by a Josephson junction (purple), acts as a coupler between neighboring resonators. The coupling is fully tunable through the external flux $\Phi_{\rm c}$.

simulators for molecular vibronic spectra [16]. The additional operations required to achieve such simulations are displacement and squeezing operations, which are readily carried out using superconducting circuits [17].

Boson Sampling Hamiltonian.- Boson sampling refers to the situation in which N single photons are injected in a M-modes photonic network characterized by the unitary matrix U. Introducing the Fock number basis, *i.e.* the basis composed by states $\{|n_1, n_2, \cdots, n_M\rangle\}$ having a precise number of photons n_j in each mode $j = 1, 2, \cdots, M$, we can write the input and output state as

$$|\psi_{\rm in}\rangle = |1_1, \cdots, 1_N, 0_{N+1}, \cdots, 0_M\rangle\,, \tag{1}$$

$$|\psi_{\rm out}\rangle = \hat{R}_U\,|\psi_{\rm in}\rangle\,, \tag{2}$$

where the transformation $\hat{R}_U$ is defined through its action on the bosonic creation operators by $\hat{R}_U\, a_i^\dagger\, \hat{R}_U^\dagger = \sum_j U_{ij} a_j^\dagger$. Aaronson and Arkhipov showed that sampling from the photon-number output distribution $P(n_1, n_2, \cdots, n_M) = |\langle n_1, n_2, \cdots, n_M|\,\hat{R}_U\,|\psi_{\rm in}\rangle|^2$ is a computationally hard task, provided that the number of modes $M \geq N^2$ and that the unitary U is chosen randomly according to the Haar measure [1].

Since any linear optical network can be constructed with phase shifters (ps) and beam splitters (bs) alone, $\hat{R}_U$ can also be decomposed as the sequential product of the corresponding unitary operations acting, respectively, only on one or two modes. In particular, it has been proven that any $M \times M$ unitary matrix U can be decomposed into $K = \mathcal{O}(M^2)$ optical elements connecting nearest-neighbor modes [18, 19], providing the factorization $\hat{R}_U = \hat{U}^{(K)} \cdots \hat{U}^{(1)}$. Every operation corresponds to the application of an appropriate Hamiltonian for the specific time τ_k according to $\hat{U}^{(k)} = \exp(-i\hat{H}_k\tau_k)$ and most operations involving distinct resonators can be performed simultaneously. The Hamiltonians have only two possible forms ($\hbar$=1 throughout)

$$\hat{H}_k^{\rm bs} = g_k a_{i_k}^\dagger a_{i_k+1} + {\rm H.\,c.}\,, \tag{3}$$

$$\hat{H}_k^{\rm ps} = \phi_k a_{j_k}^\dagger a_{j_k}\,, \tag{4}$$

where indexes $i_k, j_k = 1, \cdots, M$ label the resonator modes involved in the k-th operation. Once introduced in the operator $\hat{U}^{(k)}$, the quantities $g_k\tau_k$ and $\phi_k\tau_k$ define, respectively, the beam splitter reflectivity and phase shift associated to the k-th optical element. By applying these building-block operations sequentially, one realizes the complete boson sampling unitary $\hat{R}_U$. This procedure offers the possibility of implementing boson sampling in any platform capable of generating the above Hamiltonians. In particular, superconducting circuits associate an extraordinary level of control to the required interactions.

Boson sampling with superconducting circuits.- Boson sampling consists of three fundamental steps: i) initial single-photon state preparation, ii) implementation of the random unitary $\hat{R}_U$ and iii) single-photon detection. Here, we describe the specific circuit design to implement all the necessary operations with microwave photons.

Our proposal consists of a series of high-quality-factor (high-Q) superconducting storage resonators which are coupled to each other by a tunable interaction that can effectively be switched on and off. These resonators are used for storage of the photons that will be processed to carry out the boson sampling algorithm. At the same time, each storage resonator is also coupled to a superconducting qubit to perform the crucial operations required by a boson sampling device (see Figure 1). While this proposal is largely qubit independent, we have chosen to illustrate it adopting the X-mon qubit [20] since its cross-shaped design allows for both transverse and longitudinal rotations without acting on the resonator themselves. Finally, the design requires the X-mon qubit to be also coupled to a low-Q (measurement) resonator, which will be used to perform quantum non-demolition detection of the photons stored in the storage resonator. We now discuss how to implement all the fundamental operations in superconducting setups:

I) Initial state preparation: We initialize the qubits, initially far detuned in energy from the storage resonator frequency, in the ground state $|g\rangle$. Then, we coherently drive the first N X-mon qubits through their XY ports to implement a π-pulse that brings the qubits to the excited state $|e\rangle$. This single qubit operation can be done with extremely high fidelity, of around 99.92% as recently reported in a similar system [12, 21]. By tuning the X-mon frequency through the Z qubit control line, we bring the qubits on resonance with the storage resonators for a time t, activating a Jaynes-Cummings interaction of the

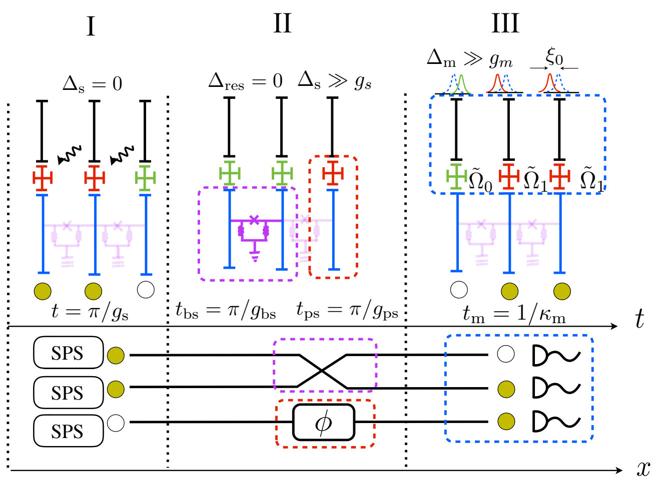

FIG. 2. Pictorial description of microwave boson sampling in a three-mode device, and comparison with its linear optics counterpart. In the optical network photons generated from a non-deterministic single photon source (SPS) travel from left to right, passing through the three fundamental steps of boson sampling I) state preparation, II) unitary dynamics and III) detection. The corresponding operations in circuit QED are illustrated in the panel above, where the color code indicates which interaction is currently active. Qubits are depicted in red if in the excited state $|e\rangle$, and green if in the ground state $|g\rangle$. Fock state $|0\rangle$ ($|1\rangle$) is shown as an empty (full) circle. Purple ring coupler are disconnected when faded. The protocol is summarized in Table I.

form

$$H_{\rm JC} = \omega_{\rm s} a^\dagger a + \frac{\Omega}{2}\sigma_z + g_s(\sigma^+ a + \sigma^- a^\dagger)\,, \qquad (5)$$

where Ω is the qubit frequency, ω_s the storage resonator frequency and $g_{\rm s}$ is the coupling constant. Applying this interaction for a time $t = \pi/g_{\rm s}$ moves the qubit excitation onto the storage resonator $|e\rangle \otimes |0\rangle \to |g\rangle \otimes |1\rangle$, creating a single-photon Fock state on the storage resonator. This operation can be performed deterministically and with high efficiency, as shown in [17]. Interestingly enough, we are not limited to the generation of single-photon states. More complicated states, such as higher-number Fock states [22] and Gaussian states [23], can also be prepared. As we will discuss later on, this would allow the implementation of boson sampling with modified input states in the form required by the quantum simulations of molecular spectroscopy [16].

II) Unitary operation: Beam splitter operations of the form (3) can be simply carried out by bringing two transmission line resonators together. In the confluence of their center conductors, evanescent waves couple the two resonators allowing the photons to tunnel between them. However, their coupling is determined by the fixed geometric arrangement of the resonators, resulting in a static coupling $g_{\rm bs}$ that can not be switched off. In order to make the coupling switchable, different schemes have been proposed theoretically [24–26] and implemented experimentally [13–15]. All these proposals are based on superconducting rings acting as tunable couplers (cf. Fig. 1a). Switchability relies on a controlled quantum interference between the resonator wavefunctions, that either adds them up or cancels each other out, depending on a control parameter, namely the external magnetic flux $\Phi_{\rm c}$ threading the superconducting ring [25]. These tunable interactions have been realized both as qubit-qubit [13] and as resonator-resonator couplers [14, 15], reporting on-off interaction ratios of about 10^4. Moreover, the switching operation is very fast and takes only a fraction of a nanosecond, that is to say a time scale much faster than the resonator dynamics.

Phase-shifting operations can be implemented by bringing the qubit off-resonance with the storage resonator, in the so-called dispersive regime where $\Delta_s = \Omega - \omega_s \gg g_s$. Under this condition, the qubit induces a state-dependent pull of the resonator frequency of the form

$$H_{\rm dis} = \left(\omega_s - \frac{g_s^2}{\Delta_s}\sigma_z\right) a^\dagger a + \frac{1}{2}\left(\Omega - \frac{g_s^2}{\Delta_s}\right)\sigma_z, \qquad (6)$$

where the effective resonator frequency includes contribution from $\phi = g_s^2/\Delta_s \times \langle\sigma_z\rangle$. As a consequence, the phase accumulated by each photon in the resonator depends on the qubit state, being proportional to $\langle\sigma_z\rangle = \pm 1$ for the excited and ground state, respectively. Assuming that every qubit is in the ground state $|g\rangle$, and equally detuned with respect to its storage resonator, there is no relative frequency shift between resonators. However, relative phases between resonators can be arbitrarily created simply by flipping the corresponding qubit to its excited state $|e\rangle$ and introducing a frequency modification equals to 2ϕ. Thus, applying a dispersive interaction of the form (6) to a desired qubit-resonator pair for times $t_{\rm ps} \in [0, \pi/\phi]$ one can introduce arbitrary *relative* phase-shifts between any pair of adjacent storage resonators.

III) Readout: The first mechanism we envision consists of mapping the storage resonator state back to the qubits, by inverting the state preparation procedure. This mechanism is supposed to perfectly distinguish between an empty resonator and a resonator occupied by a single microwave photon, as required in the original formulation of boson sampling [1]. Bringing the qubits on resonance with the storage resonators, the interaction in eq. (5) causes Rabi oscillations that swap the boson sampling resonator state $|\psi_{\rm out}\rangle$ to the qubit [17]. While two or more photons might have bunched together on the same resonator, thus preventing the transfer to the qubit state due to a photon-blockade effect [27], we can postselect this event as we would do in any linear optics implementation. With the aid of a second, low-Q resonator, we perform a quantum non-demolition detection of the qubit state (see Figure 1). Measuring the transmission of the measurement resonator, we detect with large fidelity whether the qubits are in the ground or excited state, and hence the photon state in the storage resonators [28].

	Step I: Initial state preparation	Step II: Unitary operator		Step III: Measurement protocol
		Beam splitter	Phase sifting	
Physical system	qubit-storage resonator	resonator-resonator	qubit- storage resonator	qubit- measurement resonator
Hamiltonian	Jaynes-Cummings	beam-splitting	dispersive	dispersive
Relevant parameters	$\Delta_{\rm s} = 0$, $t = \pi/g_{\rm s}$	$\Delta_{\rm res} = 0$, $t_{\rm bs} = \pi/g_{\rm bs}$	$\phi = g_s^2/\Delta_{\rm s}$, $t_{\rm ps} = \pi/\phi$	$\xi_0 = g_m^2/\Delta_{\rm m}$, $t_{\rm m} = 1/\kappa_m$
Figures of merit	$g_{\rm s}/2\pi \simeq 150$MHz	$g_{\rm bs}/2\pi \simeq 30$MHz	$\phi \simeq 20$MHz, $\kappa_s/2\pi = 1$ KHz	$\xi_0 = 30$MHz, $\kappa_m/2\pi = 20$MHz

TABLE I. Summary of the microwave boson sampling implementation. For each step of the protocol (columns), we display the key physical systems involved in that step, the Hamiltonian ruling the system dynamics, as well as the relevant parameters and their figures of merit. In step I, qubit and storage resonator interact of resonance via Jaynes-Cummings Hamiltonian. In step II, storage resonators are coupled on resonance via beam-splitter interaction for the purposes of beam splitting operations, while an off resonance, dispersive interaction with the qubit implements relative phase shifts. In step III, off resonance dispersive interaction, this time with the measurement resonator, is used for quantum non-demolition detection.

A second readout mechanism works as a *high-efficient quantum non-demolition* photon counter. The measurement mechanism is based on qubit-photon logic gates [29]. Within the dispersive regime, where the qubit is detuned by an amount Δ_m, the effective qubit frequency is lifted due to the photons in the storage resonator according to $\tilde{\Omega}_n = \Omega + \xi_n$, where the dispersive shift $\xi_n = (2n+1)g_{\rm s}^2/\Delta$, depends on the number of photons n on the storage the resonator. Then, by sending coherent microwave signals at the different frequencies of the qubit $\tilde{\Omega}_n$, we perform a π-rotation on the qubit, contingent on the storage resonator state $|n\rangle$: when the driving microwave hits the qubit at its resonant frequency, we flip the qubit state $|g\rangle \to |e\rangle$, which will, in turn, create a displacement of the measurement resonator frequency [28]. By tracking the transmission on the measurement resonator, we determine the number of photons n in the storage resonator in at most n trials. As far as boson sampling is concerned this would normally correspond to one or two attempts to measure the resonator. Each readout can be performed with efficiency of about 90% [29] and, since the measurement is non-demolition, one can repeat the measurement many times to exponentially reduce the probability of failure. The latter readout scheme represents a remarkable feature of our microwave setup that is absent, in its deterministic form, in linear optical setups. Figure 2 illustrates a pictorial comparison of a three-mode boson sampling implementation with superconducting circuits and the original linear optical network. A summary of the whole microwave boson sampling protocol can be found in Table I, where we present the most relevant parameters together with their experimental benchmarks.

Generalized boson sampling in circuit QED. Consider the very same device presented to tune the resonator-resonator couplings. The specific form of the interaction in eq. (3) is obtained in the rotating wave approximation starting from the more accurate form $H_{\rm int} = g_k(\Phi_{\rm c})(a_k^\dagger + a_k)(a_{k+1}^\dagger + a_{k+1})$. As detailed in [25], when the external magnetic flux through the coupler $\Phi_{\rm c}$ oscillates at the appropriate frequency $\omega_{\rm c} = \omega_k + \omega_{k+1}$, the interaction effectively produces two-mode squeezing in the frame rotating at the coupler frequency. Simultaneously, a displacement operation can be straightforwardly introduced by simply driving the storage resonator itself. The combined action of displacement and squeezing is interpreted as the required state preparation step of modified boson sampling setups [3, 16]. The other essential requirement is the ability of determining the parity or counting the number of photons in a resonator. In essence, we exploit the nearby qubit to check a single occupation number of the storage resonator at a time, effectively implementing a quantum non-demolition photon counter. By virtue of the suggested protocol, our proposal constitutes, to the best of our knowledge, the first scalable implementation of any practical application of boson sampling.

Discussion and Conclusion.- To address the feasibility of our proposal, we have to understand how the requirements on the single operation affect the overall scalability. Loading and measuring the resonator is performed only once per run, while a typical operation consisting of a beam splitter followed by a phase shifter requires a time $(t_{\rm bs} + t_{\rm ps}) \simeq 0.3\mu s$ (see values reported in Table 1). This time has to be compared with the storage resonator lifetime, which would probably be the limiting factor to run a successful experiment. High finesse coplanar waveguides resonators with quality factors above one million have been reported [30, 31], yielding cavity decay rates $\kappa \simeq 2\pi \times 1$ KHz corresponding to a cavity lifetime $t_\kappa = 50\mu s$. Thus, one could implement a total number of operations $t_\kappa/(t_{\rm bs} + t_{\rm ps}) \simeq 500$ before the photons are lost. Since boson sampling is believed to be hard for $N \sim \sqrt{M}$, we can successfully manipulate ~ 20 photons.

The above analysis provides a preliminary estimate based solely on the ratio between the coherence time of superconducting devices and the time for a single operation. A more careful analysis should involve the fidelity of those operations as well. We discuss the three steps separately. Firstly, loading a single photon can be done with an overall fidelity $\mathcal{F}_I = \mathcal{F}_1 \times \mathcal{F}_2$, where $\mathcal{F}_1$ is the fidelity of the qubit π-pulse and $\mathcal{F}_2$ the fidelity of the qubit-resonator swap. Experimentally demonstrated performance such as $\mathcal{F}_1 = 99.9\%$, and $\mathcal{F}_2 = 99.4\%$ [32] yields a success probability $P = (\mathcal{F}_I)^N \simeq 0.87$ for $N = 20$. Secondly, Leverrier et al. [33] showed that the average fidelity on each beam splitter and phase shifter operation should scale as $\mathcal{F}_{II} = 1 - \mathcal{O}(1/N^2)$ in order to implement a unitary U that provides a classically-hard output probability distribution. Recently reported qubit-resonator swap fidelities $\mathcal{F}_{\rm swap} = 99.4\%$ that are arguably very close to the required $\mathcal{F}_{II} \approx 99.75\%$ for a successful 20-photon experiment. Finally, the total fidelity for the measurement step is given by $\mathcal{F}_{III} = \mathcal{F}_2 \times \mathcal{F}_3$, where $\mathcal{F}_3$ is the qubit read-out fidelity. Assuming the resonators are not thermally populated, and using $\mathcal{F}_3 = 99\%$ one gets a success probability $P = \mathcal{F}_{III}^N \simeq 0.80$. These numbers are well above the thresholds reported in other boson sampling platforms, and show that the first post-classical computation is within experimental reach with today's technology using superconducting circuits.

In conclusion, we propose a realistic architecture for scalable boson sampling with superconducting circuits, which allows for deterministic state preparation and photon counting. Moreover, nonlinear operations such as squeezing are now available to realize the first practical application of boson sampling in the context of molecular spectroscopy.

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