Deterministic realization of a universal quantum gate in a single scattering process

F. Ciccarello[1,3], D. E. Browne[2], L. C. Kwek[3], H. Schomerus[4], M. Zarcone[5], and S. Bose[2]
[1]*Scuola Normale Superiore, Piazza dei Cavalieri, 7, I-56126 Pisa, Italy*
[2]*Department of Physics and Astronomy, University College London, Gower Street, London WC1E 6BT, United Kingdom*
[3]*Centre for Quantum Technologies, National University of Singapore, 3 Science Drive 2, Singapore 117543*
[4]*Department of Physics, Lancaster University, Lancaster, LA1 4YB, United Kingdom*
[5] *CNISM and Dipartimento di Fisica, Università degli Studi di Palermo, Viale delle Scienze, Ed. 18, I-90128 Palermo, Italy*

We show that a flying particle, such as an electron or a photon, scattering along a one-dimensional waveguide from a pair of static spin-1/2 centers, such as quantum dots or superconducting qubits, can implement a CZ gate (universal for quantum computation) between them. This occurs deterministically in a single scattering event, hence with no need for any post-selection or iteration, and without demanding the flying particle to bear any internal spin. We show that an easily matched hard-wall boundary condition along with the elastic nature of the process are key to such performances.

PACS numbers: 03.67.Lx, 03.67.Hk, 42.50.-p

Interfacing static qubits [1] mediated by flying particles is a prominent paradigm in the quest for efficient ways to implement quantum information processing (QIP) [2, 3]. As a major motivation, this is the only way to *jointly* address quantum registers located *far* from each other, thus featuring no direct mutual interaction (this is usually sought to favor local addressing). Within this general framework, over the past few years a research line has thrived around the idea that the crosstalk between the static qubits can be mediated by particles *scattering* from them.

Yet, all of such strategies unavoidably face an inherent major drawback. For a given quantum task to be efficiently accomplished, the link between the static objects should occur by means of the local interaction of each static object with a *quantum* flying bus. Namely, this should feature inherently quantum motional and/or internal (pseudo) spin degrees of freedom (DsOF). To do so, however, the coupling between the flying and static particles will in general entangle them so as to bring about decoherence affecting the DsOF of the static objects. Owing to such effect, the attainment of satisfactory figures of merit thus demand for further actions to complement the above interaction. Such actions typically comprise iterated injection of the flying particles and post-measurements over their DsOF [2–4]. While such conditioning of the bus dynamics generally enhances the performances, it usually comes at the cost of making the process probabilistic and, moreover, may demanding in practice (for instance, the post-selection of mobile-electron spins in semiconducting media [5]).

Moreover, as far as *scattering* scenarios are concerned, there appear intrinsic hindrances in the accomplishment of certain tasks such as the implementation of a two-qubit quantum gate (TQG), typically the most challenging as well as essential process in most QIP architectures especially when allowing for universal quantum computation (QC). To see this, assume that we need to realize a TQG between a flying qubit and a scattering center (SC) endowed with spin in a one-dimensional (1D) waveguide [6]. Even if one conditions the dynamics to either the reflection or transmission channel, the resulting process lacks unitarity, i.e. a paramount prerequisite for a TQG, unless quite specific regimes of parameters and, importantly, interaction models are addressed [6]. Analogous considerations *a fortiori* hold when many scattering centers are present and a gate involving their DsOF is sought, which will be our focus in this work. In general, scattering-based scenarios thus appear as adverse arenas to perform quantum *algorithms*, which arguably explains why mere entanglement generation was almost exclusively investigated to date [4].

Despite the above, scattering-based implementations are attractive because of the low demand for control entailed in scattering process. One normally just needs to set the wave vectors of the itinerant buses and wait for the collision to occur, thus bypassing any tuning in the interaction-time, (which is usually a significant noise source). Further benefits such as the resilience to relevant detrimental factors including static disorder, phase noise and imperfect particle-wave-vector setting [4] have been shown. Except for the attempt in Ref. [6] such advantages have so far been harnessed solely for mere entanglement generation [4].

Here, we discover a simple strategy for the realization of quantum gates between static qubits through a particle scattering from them. The injection of the latter, which we do not demand to bear any internal (pseudo) spin, followed by its multiple scattering from the SCs suffice to *deterministically* achieve the gate in one shot. To this goal, *neither post-selection of any kind nor repeated sending of the flying mediators* are required. Thus, besides unveiling an unexpected suitability of scattering-based methods [4] to achieve unitary operations, our work sets a milestone within the distributed-QIP context [2, 3]. For two static qubits a *universal* CZ gate naturally arises, which shows the effectiveness of our scheme.

Central idea. The idea behind our work is quite simple and can be illustrated as follows. Consider a monochromatic spinless particle f of wave vector k propagating along a 1D wire that impinges on an array of SCs [see Fig. 1(a)]. Once multiple scattering has occurred and assuming, importantly, that the process is elastic, f can only be found either reflected or transmitted with corresponding wave vectors $-k$ and k, respectively, as shown in Fig. 2(b). Let $\{|\nu\rangle\}$ be a basis of the SCs' Hilbert space and $|\mu\rangle$ one of its elements, whereas $|\pm k\rangle$ are momentum eigenstates of f. Let $|\Psi_{in}\rangle = |k\rangle|\mu\rangle$ be the overall system's state

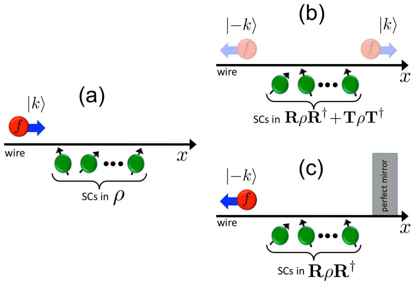

FIG. 1. (Color online) Scheme working principle. (a) f impinges on a set of SCs in the state ρ. (b) After the interaction, f is either reflected or transmitted. Concomittantly, the SCs undergo a non-unitary quantum map. (c) With a perfect mirror beyond the centers, f can be only back reflected and the unitary $\hat{R}$ is applied to the SCs.

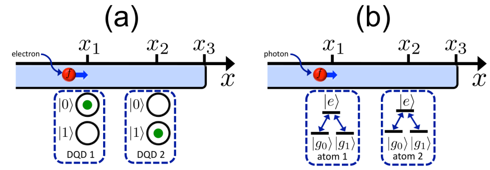

FIG. 2. (Color online) Setups implementing our scheme. (a) An electron along a quantum wire with DQDs. (b) A photon along a waveguide with Λ-type atom-like systems. The right end of the wire/waveguide behaves as a perfect mirror.

prior to scattering. As f is scattered off the final state reads

$$|\Psi_f\rangle = |-k\rangle \sum_\nu r_{\nu\mu}|\nu\rangle + |k\rangle \sum_\nu t_{\nu\mu}|\nu\rangle \ , \qquad (1)$$

where $r_{\nu\mu}$ ($t_{\nu\mu}$) is a reflection (transmission) probability amplitude corresponding to the initial and final centers' states $|\mu\rangle$ and $|\nu\rangle$, respectively. Defining a reflection (transmission) operator $\hat{R}$ ($\hat{T}$) in the Hilbert space of the SCs such that $\langle\nu|\hat{R}|\mu\rangle = \mathbf{R}_{\nu\mu} = r_{\nu\mu}$ ($\langle\mu|\hat{T}|\nu\rangle = \mathbf{T}_{\nu\mu} = t_{\nu\mu}$) Eq. (1) can be arranged as $|\Psi_f\rangle = |{-k}\rangle\hat{R}|\mu\rangle + |k\rangle\hat{T}|\mu\rangle$. Tracing over f, the final density operator of the centers is thus obtained as $\hat{R}|\mu\rangle\langle\mu|\hat{R}^\dagger + \hat{T}|\mu\rangle\langle\mu|\hat{T}^\dagger$. This immediately yields that when the centers are initially in an arbitrary state described by the density matrix ρ (in general mixed) their final state is given by

$$\rho' = \mathbf{R}\rho\mathbf{R}^\dagger + \mathbf{T}\rho\mathbf{T}^\dagger \ . \qquad (2)$$

Due to the normalization condition $\sum_\nu(|r_{\nu\mu}|^2 + |t_{\nu\mu}|^2) = 1$ for any μ the above operators fulfill $\hat{R}\hat{R}^\dagger + \hat{T}\hat{T}^\dagger = \mathbb{1}$, where $\mathbb{1}$ is the identity operator of the SCs' Hilbert space. We would like the scattering process to implement a multi-qubit gate, which is *unitary*, between the SCs. In general, this is not the case as is evident from Eq. (2) showing that the SCs undergo instead a quantum map [1] comprising reflection and transmission channels.

Such hindrance can be got around in a very natural fashion by simply inserting beyond the centers a *perfect mirror* [see Fig. 1(c)]. As this introduces a hard-wall boundary condition (BC) preventing f from trespassing the right end, we in fact eliminate the transmission channel so that it is thus fully *suppressed*. Hence, $t_{\nu\mu} \equiv 0$ and (2) reduces to

$$\rho' = \mathbf{R}\rho\mathbf{R}^\dagger \ ,$$

where now $\hat{R}\hat{R}^\dagger = \hat{R}^\dagger\hat{R} = \mathbb{1}$, i.e. *in the presence of the mirror* $\hat{R}$ *becomes a unitary gate*. Thereby, such gate is deterministically implemented whenever f scatters from the centers. Note that this holds regardless of the specific scattering potential. Rather, this affects only the *type* of achieved gate.

Having illustrated how a hard-wall BC guarantees the process unitarity, it is now natural to wonder whether there exist elastic one-channel scattering processes allowing for multi-qubit gates that are *universal* for QC [1]. We will show that this is indeed the case. To this aim, we first focus on a simple paradigmatic setting, dubbed *setup A*, comprising two spin-1/2 centers, i.e. two qubits [1], each coupled to a massive particle embodying f. We identify a regime such that a CZ gate, universal for QC [1], naturally arises. We next address *setup B*, which comprises multi-level atom-like systems and photons and is within experimental reach in a manifold of scenarios [7–10]. We show that it can work as an effective emulator of setup A in that it is endowed with all the essential features enabling the CZ gate.

Setup A. This setting, sketched in Fig. 2(a), comprises two identical spin-1/2 scattering centers, 1 and 2, lying along a semi-infinite wire along the x-axis at $x = x_1$ and $x = x_2$, respectively, each coupled to a scattering particle f of mass m. The wire ends at $x = x_3$. Let $\{|0\rangle_i, |1\rangle_i\}$ be an orthonormal basis for the ith center ($i = 1, 2$). In practice, one can consider a semiconducting quantum wire or a carbon nanotube [11] where an electron populating the lowest subband can undergo scattering from two double quantum dots (DQDs) [12] to which it is electrostatically coupled. As shown in Fig. 2(a), each single-electron DQD is in $|0\rangle$ ($|1\rangle$) if the upper (lower) dot is occupied, hence implementing an effective qubit [12] (we assume negligible tunneling between the upper and lower dots during the passage of f). Also, the electrostatic coupling is negligible for state $|1\rangle$. The Hamiltonian is thus modeled as (we set $\hbar = 1$ throughout)

$$\hat{H} = \frac{\hat{p}^2}{2m} + \Gamma\delta(x - x_1)\,|0\rangle_1\langle 0| + \Gamma\delta(x - x_2)\,|0\rangle_2\langle 0| \ , \qquad (3)$$

where $\hat{p} = -i\,d/dx$ is the momentum operator of f, while Γ is the height of each contact potential scattering centered at $x = x_i$ ($i = 1, 2$). Note that the scattering potential in (3) is *dispersive* because it cannot induce either 1 or 2 to flip between $|0\rangle$ and $|1\rangle$. Upon scattering, each initial SCs' state $|\alpha_1\alpha_2\rangle_{12}$ ($\alpha_i = 0, 1$) simply but crucially picks up its own phase shift. It is trivially checked in the light of (3) that for a given state labeled by $\boldsymbol{\alpha} \equiv \{\alpha_1, \alpha_2\}$ $\hat{H}$ takes the effective form $\hat{H}_{\boldsymbol{\alpha}} = \frac{\hat{p}^2}{2m} + \sum_{i=1,2}\Gamma\delta_{\alpha_i 0}\,\delta(x - x_i)$. The problem thus reduces to a particle scattering from spin-less potentials. We label with $r_{\boldsymbol{\alpha}}$ the f's reflection probability amplitude corresponding to $\hat{H}_{\boldsymbol{\alpha}}$, where the subscript here specifies both the initial and final centers' state, which coincide owing to the dispersive interaction. For the same reason, in the computational basis $\mathcal{B} = \{|00\rangle_{12}, |01\rangle_{12}, |10\rangle_{12}, |11\rangle_{12}\}$ $\mathbf{R}$ necessarily has the diagonal form $\mathbf{R}=\mathrm{diag}(\mathrm{r}_{00},\mathrm{r}_{01},\mathrm{r}_{10},\mathrm{r}_{11})$.

Adopting a standard procedure, in order to derive $r_{\boldsymbol{\alpha}}$ we assume that f is left-incoming and seek the stationary state $\Psi_{\boldsymbol{\alpha}}(x)$ fulfilling $\hat{H}_{\boldsymbol{\alpha}}\Psi_{\boldsymbol{\alpha}}(x) = k^2/(2m)\Psi_{\boldsymbol{\alpha}}(x)$ of the form $\Psi_{\boldsymbol{\alpha}}(x) = \Psi_{\boldsymbol{\alpha}+}(x)+\Psi_{\boldsymbol{\alpha}-}(x)$ with (to simplify the notation we drop the dependance on $\boldsymbol{\alpha}$ whenever unnecessary)

$$\Psi_+(x)=\{\theta[-x)+a_1\,[\theta(x)-\theta(x-x_2)]+a_2\theta(x-x_2)\}\,e^{ikx}\,, \quad (4)$$
$$\Psi_-(x)=\{r\,\theta(-x)+b_1\,[\theta(x)-\theta(x-x_2)]+b_2\theta(x-x_2)\}\,e^{-ikx}\,, (5)$$

where $k>0$ is the wave vector modulus, $\theta(x)$ is the Heaviside step function and we have set $x_1=0$. In the light of Eqs. (4) and (5), $\Psi_+(x)$ [$\Psi_-(x)$] represents the right-propagating (left-propagating) part of $\Psi(x)$. Note that $\Psi(x)$ is specified by the five coefficients $\{r,a_1,b_1,a_2,b_2\}$, which do depend on $\boldsymbol{\alpha}$. These coefficients can be found by requiring that $\Psi(x)$ and its derivative with respect to x $\Psi'(x)$ match the five BCs

$$\Psi(x_i^-) = \Psi(x_i^+) \quad (i=1,2)\,, \qquad \Psi(x_3) = 0, \qquad (6)$$
$$\Delta\Psi'|_{x_i} = 2m\,\Gamma\delta_{\alpha_i 0}\,\Psi(x_i) \quad (i=1,2)\,. \qquad (7)$$

Eqs. (6) ensures the matching of the wave function at the centers' locations (first two) and the hard-wall BC owing to the end of the wire at $x=x_3$ (latter equation). Eqs. (7) are standard constraints on the discontinuity of $\Psi'(x)$ at the centers' positions $\Delta\Psi'|_{x_i}=\Psi'(x_i^+)-\Psi'(x_i^-)$ due to the δ-potentials in (3) [they are derived by integrating the Schrödinger equation (SE) corresponding to $\hat{H}_{\boldsymbol{\alpha}}$ over an infinitesimal range across $x=x_1$ and $x=x_2$]. Solving the linear system (6)-(7) we end up with

$$r_{\boldsymbol{\alpha}}=-\exp\{2i\arg[\exp(-ikx_{31})-2\gamma\delta_{\alpha_2 0}[\cos kx_{21} -(i+2\gamma\delta_{\alpha_1 0})\sin kx_{21}]\sin kx_{32}]-2\gamma\delta_{\alpha_1 0}\sin kx_{31}]\}\,, \quad (8)$$

where we have set $x_{ij}=x_i-x_j$ and $\gamma=m\Gamma/k$. Evidently, as expected, the above yields $|r_{\boldsymbol{\alpha}}|=1$ regardless of all the parameters and $\boldsymbol{\alpha}$, namely f is reflected back with certainty. Let us now focus on the regime defined by

$$kx_{21}=n\pi, \quad kx_{32}=(n'+1/2)\,\pi \quad \gamma\gg 1\,, \qquad (9)$$

where n and n' are arbitrary integers. Replacing (9) in (8) and using that $kx_{31}=kx_{32}+kx_{21}$, we end up with

$$r_{00}=r_{01}=r_{10}=-r_{11}=-1\,, \qquad (10)$$

which yields the gate matrix as $\mathbf{R}=\mathrm{diag}(1,1,1,-1)$ (up to an irrelevant global phase factor), i.e. the well-known CZ gate [1]. This proves that a single scattering process can implement a universal TQG. Intuitively, unlike Ref. [6] here the gate unitarity is secured by the geometry, which leaves the physical parameters free to be tuned in a way that a CZ is matched.

Setup B. In this setting, sketched in Fig. 2(b), f is a photon propagating along a 1D waveguide, whose geometry is analogous to the one of setup A, and scattering from two three-level atom-like systems. Each "atom" $i=1,2$ has a Λ-type energy-level configuration consisting of a twofold-degenerate ground doublet spanned by states $\{|g_0\rangle,|g_1\rangle\}$ and an excited state $|e\rangle$ [see Fig. 2(b)]. Hence, a two-photon Raman transition between $|g_0\rangle$ and $|g_1\rangle$ can occur through absorption and re-emission of a scattering photon. Assuming a linear photon dispersion law $\omega=\upsilon k$ with ω the photon energy and υ the group velocity, the free-field Hamiltonian in the waveguide can be written as [13, 14] $\hat{H}_f=-i\sum_{d=\pm}\int dx\,\upsilon_d\,\hat{c}_d^\dagger(x)\partial_x\hat{c}_d(x)$ with $\upsilon_+=-\upsilon_-=\upsilon$ and $\hat{c}_+^\dagger(x)$ [$\hat{c}_-^\dagger(x)$] the bosonic operator creating a right (left) propagating photon at x. The free atomic Hamiltonian reads $H_a=\omega_0\sum_{i=1,2}|e\rangle_i\langle e|$, where ω_0 is the energy gap between $|e\rangle$ and the ground doublet. As for the interaction between the field and the ith atom, this is modeled as [14] $\hat{H}_{fi}=J\int dx\,\delta(x-x_i)\,[\hat{c}(x)\hat{S}_i^\dagger+\mathrm{H.c.}]$ (under the usual rotating-wave approximation), where the bosonic operator $\hat{c}(x)=\hat{c}_+(x)+\hat{c}_-(x)$ annihilates a photon at x regardless of its propagation direction, J is the rate associated with each transition $|g_j\rangle_i\leftrightarrow|e\rangle_i$ ($\forall j=0,1$) and $\hat{S}_i^\dagger=\sum_j|e\rangle_i\langle g_j|$. The full Hamiltonian thus reads $\hat{H}=\hat{H}_f+\hat{H}_a+\hat{H}_{f1}+\hat{H}_{f2}$. In virtue of the system's symmetries, it is now convenient to use as basis of each ground doublet the symmetric and antisymmetric combinations of the two ground states $|\phi^\pm\rangle_i=(|g_0\rangle_i\pm|g_1\rangle_i)/\sqrt{2}$. As $|\phi^-\rangle_i$ is a dark state [14], i.e. $\hat{S}_i^\dagger|\phi^-\rangle_i=0$, the atomic raising operator takes the effective form $\hat{S}_i^\dagger\equiv|e\rangle_i\langle\phi^+|$. Thus, unlike $\{|g_0\rangle,|g_1\rangle\}$, the Raman process does not couple $|\phi^+\rangle$ and $|\phi^-\rangle$. It should be clear now that by taking $|0\rangle=|\phi^+\rangle$ and $|1\rangle=|\phi^-\rangle$ for each atom, as long as these are initially in the ground doublet, setup B in fact possesses all the key features of A. Indeed, if the ith atom is in $|1\rangle_i=|\phi^-\rangle_i$ the corresponding potential $\hat{H}_{fi}$ vanishes. If it is initially in $|0\rangle_i$, it may undergo a second-order transition $|0\rangle_i\rightarrow|e\rangle_i\rightarrow|0\rangle_i$ so as to eventually pick up a phase shift once f is scattered off. To make rigorous such considerations, we next prove that the reflection coefficients are again, with due replacements, given by (8) as for setup A.

$\hat{H}$ conserves the total number of excitations. Thus, following a standard approach [13, 14], we search for one-excitation stationary states of the form

$$|\Psi_{\boldsymbol{\alpha}}\rangle=\sum_{d=\pm}\int dx\,\psi_{\boldsymbol{\alpha}d}(x)\hat{c}_d^\dagger(x)|\mathrm{vac}\rangle|\alpha_1\alpha_2\rangle_{12} +\varepsilon_1|\mathrm{vac}\rangle|e\alpha_2\rangle_{12}+\varepsilon_2|\mathrm{vac}\rangle|\alpha_1 e\rangle_{12}, \qquad (11)$$

where $\psi_{\boldsymbol{\alpha}\pm}(x)$ have a form analogous to Eqs. (4) and (5) thus being specified by parameters $\{r,a_1,a_2,b_1,b_2\}$, $\{\varepsilon_i\}$ are excited-state amplitudes and $|\mathrm{vac}\rangle$ is the field vacuum state. Thus, for given $\boldsymbol{\alpha}$ $|\Psi_{\boldsymbol{\alpha}}\rangle$ is specified by the 7 complex amplitudes $\{r,a_1,a_2,b_1,b_2,\varepsilon_1,\varepsilon_2\}$ and obeys the SE $\hat{H}|\Psi_{\boldsymbol{\alpha}}\rangle=\upsilon k|\Psi_{\boldsymbol{\alpha}}\rangle$. By projecting this onto $c_\pm^\dagger(x)|\mathrm{vac}\rangle|\alpha_1\alpha_2\rangle$ we obtain (we henceforth omit subscript $\boldsymbol{\alpha}$)

$$\mp\, i\upsilon\psi_\pm'(x)+J\sum_{i=1,2}\delta_{\alpha_i 0}\varepsilon_i\delta(x-x_i)=\upsilon k\,\psi_\pm(x)\,. \qquad (12)$$

Further projection of the SE onto $|\mathrm{vac}\rangle|e\alpha_2\rangle$ and $|\mathrm{vac}\rangle|\alpha_1 e\rangle$ immediately yields that for each $i=1,2$ $\varepsilon_i=J\delta_{\alpha_i 0}\psi(x_i)/(\upsilon k-\omega_0)$. Replacing these in (12) we are left with $\{\psi(x),\psi_\pm(x)\}$ only

$$\mp\, i\upsilon\psi_\pm'(x)+J^2/(\upsilon k-\omega_0)\sum_{\ell=1,2}\delta_{\alpha_\ell 0}\psi(x_\ell)\delta(x-x_\ell)=\upsilon k\,\psi_\pm(x)\,. \quad (13)$$

By subtracting now the equation for ψ_- from the ψ_+'s one we end up with $-i\psi'(x)=k[\psi_+(x)-\psi_-(x)]$, which trivially entails that $-i\Delta\psi'|_{x_\ell}=k(\Delta\psi_+|_{x_\ell}-\Delta\psi_-|_{x_\ell})$ holds as well for each $\ell=1,2$. Each $\Delta\psi_\pm|_{x_\ell}$ on the r.h.s. of the above can be evaluated by integrating (13) over an infinitesimal interval across $x=x_\ell$ ($\ell=$

1, 2), which straightforwardly yields $\Delta\psi_{\pm}|_{x_\ell} = \mp(i/\upsilon)J^2/(\upsilon k-\omega_0)\delta_{\alpha_\ell 0}\psi(x_\ell)$. We thus end up with

$$\Delta\psi'|_{x_\ell} = 2\,\frac{k}{\upsilon}\,\frac{J^2}{\upsilon k-\omega_0}\,\delta_{\alpha_\ell 0}\,\psi(x_\ell) \quad (\ell=1,2)\ . \tag{14}$$

This is identical to Eq. (7) of setup A once we set

$$\Gamma = \frac{J^2}{\upsilon k-\omega_0} \tag{15}$$

and note that, due to the parabolic dispersion law, in setup A the group velocity reads k/m and hence in Eq. (7) m can be replaced with the ratio between k and such velocity.

Given that $\psi(x)$, besides (14), must fulfill conditions analogous to Eqs. (6) due to the common geometry of setups A and B, we conclude that amplitudes $\{r_{\boldsymbol{\alpha}}\}$ for setting B are *identical* to (8) with the effective mass and potential height given by k/υ and (15), respectively. In passing, note that $|r_{\boldsymbol{\alpha}}|^2=1$ showing that if f is absorbed it will be re-emitted with certainty (each atom behaves as an effective qubit embodied by its ground doublet). It is now evident that setup B can be used as an emulator of A, thereby allowing for occurrence of the CZ gate [cf. Eq. (10)] under conditions (9). Interestingly, the requirement $\gamma=\Gamma/k\gg 1$ now in fact becomes the *resonance condition* (RC) $\upsilon k\simeq\omega_0$, i.e. the matching between the photon and atomic frequencies. This agrees with [13], where it was shown that the reflectivity of an atomic scatterer becomes unitary in this limit. Note that such RC is routinely matched in the lab through local-field tuning of the atomic frequencies, for instance. Also, although of a different nature the first two requirements in (9) are RCs either: The CZ gate thus in fact stems from a combination of RCs.

As anticipated, setup B can be experimentally implemented in several different ways, including photonic-crystal waveguides with defect cavities [7], semiconducting (diamond) nanowires with embedded QDs (nitrogen vacancies) [8], optical or hollow-core fibers interacting with atoms [9] and microwave transmission lines coupled to superconducting qubits [10]. Interestingly, within the latter scenario it was recently shown [15] that a hard-wall BC analogous to ours [cf. Fig. 2(b)] can benefit microwave photodetection.

Conclusions. We have shown a strategy to deterministically carry out multi-qubit gates between static qubits through single flying buses scattering from them. This is effective without demanding post-selection of any kind or iteration. The possibility to naturally implement a universal CZ gate has been proven in two different setups including 1D photonic waveguides coupled to atom-like qubits. We believe this work can set a significant milestone for future advancements in the area of distributed QIP as well as in the emerging field of quantum optics in 1D waveguides.

Acknowledgements. We acknowledge support from FIRB IDEAS (project RBID08B3FM), the National Research Foundation & Ministry of Education of Singapore, Leverhulme Trust, the EPSRC, the Royal Society and the Wolfson Foundation.

[1] M. A. Nielsen and I. L. Chuang, *Quantum Computation and Quantum Information* (Cambridge University Press, Cambridge, U. K., 2000).
[2] J. I. Cirac, P. Zoller, H. J. Kimble, and H. Mabuchi, Phys. Rev. Lett.**78**, 3221 (1997); D. E. Browne, M. B. Plenio, and S. F. Huelga, *ibid.* **91**, 067901 (2003); S. Bose, P. L. Knight, M. B. Plenio, and V. Vedral, *ibid.* **83**, 5158 (1999); J. I. Cirac, A. K. Ekert, S. F. Huelga, and C. Macchiavello, Phys. Rev. A **59**, 4249 (1999).
[3] Y. L. Lim, A. Beige, and L. C. Kwek, Phys. Rev. Lett. **95**, 030505 (2005); S. D. Barrett and P. Kok, Phys. Rev. A **71**, 060310 (2005); Y. L. Lim *et al.*, *ibid.* **73**, 012304 (2006).
[4] A.T. Costa, Jr., S. Bose, and Y. Omar, Phys. Rev. Lett. **96**, 230501 (2006); G.L. Giorgi and F. De Pasquale, Phys. Rev. B **74**, 153308 (2006); K. Yuasa and H. Nakazato, J. Phys. A: Math. Theor. **40**, 297 (2007); Y. Hida, H. Nakazato, K. Yuasa, and Y. Omar, Phys. Rev. A **80**, 012310 (2009); F. Ciccarello *et al.*, New J. Phys. **8**, 214 (2006); J. Phys. A: Math. Theor. **40**, 7993 (2007); F. Ciccarello, G. M. Palma, and M. Zarcone, Phys. Rev. B **75**, 205415 (2007); H. Schomerus and J.P. Robinson, New J. Phys. 9, **67** (2007); F. Ciccarello, M. Paternostro, G. M. Palma and M. Zarcone, Phys. Rev. B **80**, 165313 (2009); F. Ciccarello, M. Paternostro, M. S. Kim, and G. M. Palma, Phys. Rev. Lett. **100**, 150501 (2008); F. Ciccarello, M. Paternostro, G. M. Palma, and M. Zarcone, New J. Phys. **11**, 113053 (2009); K. Yuasa, J. Phys. A **43**, 095304 (2010); K. Yuasa, D. Burgarth, V. Giovannetti, and H. Nakazato, New J. Phys. **11**, 123027 (2009); A. De Pasquale, K, Yuasa, and H. Nakazato, Phys. Rev. A **80**, 052111 (2009); F. Ciccarello, S. Bose, and M. Zarcone, *ibid.* **81**, 042318 (2010); F. Ciccarello *et al.*, *ibid.* **82**, 030302(R) (2010); F. Buscemi, P. Bordone, and A. Bertoni, New J. Phys. **13**, 013023 (2011); Y. Matsuzaki and J. H. Jefferson, arXiv:1102.3121.
[5] D. D. Awschalom, D. Loss, and N. Samarth, *Semiconductor Spintronics and Quantum Computation* (Springer, Berlin, 2002).
[6] G. Cordourier-Maruri *et al.*, Phys. Rev. A **82**, 052313 (2010).
[7] A. Faraon *et al.*, Appl. Phys. Lett.**90**, 073102 (2007).
[8] M. H. M. van Weert *et al.*, Nano Lett. **9**, 1989 (2009); J. Claudon *et al.*, Nat. Photonics **4** 174 (2010); T. M. Babinec *et al.*, Nat. Nanotechnol. **5**, 195 (2010).
[9] B. Dayan *et al.*, Science **319**, 1062 (2008); E. Vetsch *et al.*, Phys. Rev. Lett.**104**, 203603 (2010); M. Bajcsy *et al.*, *ibid.* **102**, 203902 (2009).
[10] A. Wallraff et al., Nature (London) **431**, 162 (2004); O. Astafiev *et al.* Science **327**, 840 (2010).
[11] D Gunlycke *et al.*, J. Phys.: Condens. Matter **18**, S851 (2006); S. J. Tans *et al.*, Nature (London) **386**, 474 (1997).
[12] T. Hayashi *et al.*, Phys. Rev. Lett. **91**, 226804 (2003), Phys. Rev. Lett. 95, 090502 (2005); J. Gorman, D. G. Hasko, and D. A. Williams, *ibid.* **95**, 090502 (2005).
[13] Shen and S. Fan, Opt. Lett. **30** (2005); Phys. Rev. Lett. **95**, 213001 (2005).
[14] D. Witthaut, and A. S. Sørensen, New J. Phys. **12**, 043052 (2010).
[15] B. Peropadre *et al.*, arXiv:1101.0016.