Energy-efficient entanglement generation and readout in a spin-photon interface

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We consider a quantum interface made of a spin coupled to a one-dimensional atom, and study its potential for energy efficient entanglement generation and readout. We show that quantum superpositions of zero and single photon states outperform coherent pulses of light, producing more entanglement with the same energy. Entanglement is generally distributed over the polarisation and the temporal degrees of freedom, except for quasi-monochromatic pulses whose shape is preserved by light-matter interaction. The energetic advantage provided by quantum pulses over coherent ones is maintained when information on the spin state is extracted at the classical level. The proposed schemes are robust against imperfections in state of the art semi-conducting devices.

Introduction - Mapping stationary qubits onto flying qubits, and vice-versa, is a key task for quantum computation [1], quantum internet [2], and quantum repeaters [3]. Due to their weak interaction with the environment, photons provide excellent flying qubits [4, 5]. On the other hand, solid-state spins, with their long-lasting coherent lifetime, implement high quality stationary qubits for information storage and processing. Spin-photon interfaces (SPIs) are devices where a travelling light field (the flying qubit) coherently interacts with a spin (the stationary qubit) via a quantum emitter. Due to optical selection rules, different polarizations drive different transitions of the spin-quantum emitter system, a key mechanism for spin-light entanglement generation, quantum information mapping, and spin-photon and photon-photon gates. In the past decade, SPIs have been implemented in a wide variety of architectures employing strong perpendicular magnetic fields [6–11], lifting the degeneracy among the transitions corresponding to different spin projections (Voigt configuration). If such a configuration can generate high-fidelity entanglement between a single spin and an emitted photon, the process cannot be used to copy spin states on incoming light states. Such operations rather require a degenerate configuration [12], see Fig.1(b).

Degenerate configurations are technologically highly demanding, as they forbid the use of strong magnetic fields. So far, they have been achieved in a few setups based on dark-excitons [13]. Degenerate SPIs have been theoretically studied with high-energy coherent fields [12] and single photons as input pulses [14–17]. They rely on the giant optical Faraday rotation – a non-linear mechanism which can be engineered by coupling the quantum emitter to a cavity and exploiting the spin-dependent cavity frequency. Refs. [14–17] focused on monochromatic photons, and cavity losses as the only source of imperfections.

In this Letter, we focus on a degenerate SPI based on a lossless “one-dimensional atom” (1D atom)(Fig. 1). A 1D atom features a quantum emitter efficiently coupled to a one-dimensional reservoir of electromagnetic modes. A semiconductor quantum dot inserted in a micropillar has been shown to be a nearly ideal 1D atom at low temperature [18, 19]. We first focus on the potential of the SPI for spin-light entanglement generation. Based on a collisional model [20, 21], we compute the exact expression of the spin-light entangled states and draw analytic comparisons between resonant input pulses of arbitrary statistics, energy and bandwidth. We consider coherent input pulses, and quantum superposition of zero and one photon, that have been recently proven to be experimentally producible [22]. We show that quantum superpositions outperform coherent light pulses: unlike the latter, the former can always produce maximally entangled spin-light states for arbitrary input energies between 0.5 and 1 photon.

After the entangling process, light carries information on the spin state. In a second step, we explore if the quantum advantage also characterizes the performances for spin readout. To do so, we compare the distinguishability of the spin states, when the output light pulses undergo a classical measurement. We single out a regime where quantum superpositions outperform coherent ones, and show that the signal is robust against standard imperfections afflicting state of the art semi-conducting devices. Optimal performances are obtained in the limit of nearly monochromatic pulses, whose shape is not altered by the interaction with the spin. This regime is promising for the engineering of photon-photon gates in optical quantum computing.
System and model - The SPI features a 4-level system encompassing two degenerate transitions, each respectively coupled to a one-dimensional (1D), circularly polarized field (see Fig. 1(a)). The frequencies of the field modes are denoted $\omega_k$ with lowering operators $a_{j,k}$, where $j \in \{R,L\}$ stands for right- and left-circular polarization. The field dispersion relation $k = \omega_j v_k$ where $v$ is the field group velocity and $k$ its wave number. This system can be physically implemented, by embedding the emitter in a semi-infinite waveguide closed by a perfectly reflecting mirror (see Fig.1(b)), or equivalently, by weakly coupling the quantum emitter to an asymmetric directional cavity [14, 23]. This geometry allows us to only consider positive values of $k$. The bare field Hamiltonian reads $H_{0,f} = \hbar \sum_{j=R,L} \left[ \sum_{k=0}^{\infty} \omega_k a_{j,k}^\dagger a_{j,k} \right]$. The 4-level system is composed of two ground states, $|\uparrow\rangle, |\downarrow\rangle$, with spin projections respectively $\pm \hbar/2$, and zero energy; and two excited states, $|\uparrow\rangle, |\downarrow\rangle$, with spin projections respectively $\pm \hbar/2$ and energy $\hbar \omega_0$. Its bare Hamiltonian reads $H_{0,q} = \hbar \omega_0 \sum_{j=R,L} \sigma_j^x \sigma_j^t$, where we defined the lowering operators $\sigma_\uparrow = |\downarrow\rangle \langle \uparrow|$, and $\sigma_R = |\uparrow\rangle \langle \downarrow|$ and $\sigma_L = |\downarrow\rangle \langle \uparrow|$. Due to the conservation of the angular momentum, the transition $|\downarrow\rangle \to |\uparrow\rangle$ (resp. $|\uparrow\rangle \to |\downarrow\rangle$) is excited by left (right) circularly polarized photons.

The dynamics is solved in the interaction picture with respect to $H_0 = H_{0,f} + H_{0,q}$, yielding the interaction Hamiltonian,

$$H_I(t) = i\hbar \sqrt{2} \sum_{j=R,L} \left[ \sigma_j^x(t) b_j(t) - b_j^\dagger(t) \sigma_j(t) \right].$$

We implicitly assumed that the light-matter interaction is the same for both polarizations, and weak enough that only frequency modes close to $\omega_0$ play a role (quasi-monochromatic approximation) [24]. In this regime, the rotating wave approximation is allowed [25], and the coupling $g$ can be considered uniform in frequency [26]. We defined the emitter’s dipoles $\sigma_j(t) = e^{-i\omega_0 t} \sigma_j^t$, and temporal modes annihilation operator [20, 24, 27], $b_j(t) = \sqrt{2} \sum_{k=1,2} e^{-i\omega_k t} a_{j,k}$, with $\tau$ being the density of modes in the waveguide and $\gamma = g^2 \tau$ being spontaneous emission rate of each transition. The operator $b_j(t)$ obeys the bosonic algebra $[b_j(t), b_{j'}(s)] = \delta_{jj'} \delta(t-s)$, with $j, j' \in \{R,L\}$. In this scheme, states solely evolve under light-matter interaction: the input (resp. output) field corresponds to the initial light state at $t = 0$ (resp. the scattered light state at $t \to +\infty$).

Noteworthy, product states of the form $|\uparrow\rangle / |\downarrow\rangle_{\text{L(R)}}$ (where $|\uparrow\rangle_{\text{L(R)}}$ stands for any left/right polarized state of light) define subspaces which remain uncoupled by the dynamics, each subspace evolving like a spinless 1D atom. $|\uparrow\rangle, |\downarrow\rangle_{\text{L(R)}}$ are preserved by the scattering process. Conversely, input pulses $|\psi_{\text{in}}\rangle_{\text{L(R)}}$ (resp.$|\psi_{\text{in}}\rangle_{\text{R(L)}}$) interact with the quantum emitter in $|\uparrow\rangle$ (resp.$|\downarrow\rangle$). Resonant, low energy pulses define the linear regime where the scattering induces a $\pi$ phase shift [28] on the spin-light state. The phase shift gets reduced out of the linear regime, i.e. when the transition gets saturated by the input light. Then spontaneous emission takes place, which alters the shape of the scattered pulse. For single photons, the linear regime is equivalent to the quasi-monochromatic regime where they undergo the map:

$$|\uparrow\rangle (|\downarrow\rangle, |\text{R(L)}\rangle) \to -|\uparrow\rangle (|\downarrow\rangle, |\text{R(L)}\rangle).$$

$|\text{R(L)}\rangle$ stands for a R(L) circularly polarized single photon.

Spin-light entanglement - We first focus on the SPI potential for spin-light entanglement generation. To do so, the joint light-matter system is initially prepared in the pure product state

$$|\psi(0)\rangle = |\psi_H\rangle \otimes (|\uparrow\rangle + |\downarrow\rangle) / \sqrt{2},$$

where $|\psi_H\rangle$ stands for some horizontally (H) polarized input field. In the long-time limit, the probability that the quantum emitter is excited is negligible and the joint state reads:

$$|\psi(t \to \infty)\rangle = (|\psi_1\rangle |\uparrow\rangle + |\psi_1\rangle |\downarrow\rangle) / \sqrt{2},$$

where we dropped the explicit time-dependence in the field states to shorten the notation. Below we refer to $|\psi_1\rangle$ as the light pointer states [29]. Since the state of the joint system is pure, their overlap can be used to quantify the amount of spin-light entanglement through the so-called quantum Bhat-tacharyya coefficient (qBhat) [30, 31]:

$$B_q = |\langle \psi_1 | \psi_1 \rangle|.$$  

$B_q = 0$ implies that the pointer state measurement gives perfectly strong measurement of the spin. We first consider a resonant coherent input pulse of amplitude $a(t) = \langle \psi_1 | b_{\text{in}}(t) | \psi_1 \rangle$, with $b_{\text{in}}(t) \equiv (b_{\text{R}}(t) + b_{\text{L}}(t)) / \sqrt{2}$. The wave-function of the global system is computed using the collisional model presented in [20]: each temporal mode of the field is a unit propagating freely in the waveguide and interacting with the quantum emitter once. The unitary evolution of the joint system is decomposed as a product of independent collisions. The pointer states appearing in Eq.(3) read

$$|\psi_{\uparrow,\downarrow}^{\text{cs}}\rangle = D(\alpha_j) \left[ \sqrt{p_0} + \sqrt{p_1} \int dt f^{\dagger}(t)b_{\text{R}(\text{L})}^\dagger(t) + .. \right] |0\rangle.$$

The superscript cs refers to the coherent state and $D(\alpha_j) = e^{i\theta(\alpha_j) - \beta(\alpha_j)}$ is the displacement operator. $p_n$ stands for the probability that the quantum emitter spontaneously emits $n$ photons in the empty modes of the waveguide (see Supp. Mat. [32]). Plugging Eq. (5) into Eq. (4) yields:

$$B_q^{\text{cs}} = p_0.$$  

The physical meaning of Eq. (6) is transparent: perfect spin-light entanglement is generated, as soon as at least one spontaneous photon is emitted ($p_0 = 0$). Indeed the polarization of the photons carry complete information on the excited transition and hence, on the spin state. Figure 2 shows the value of qBhat for a coherent input pulse of amplitude
\( \alpha(t) = \sqrt{\bar{n}} e^{-\Gamma t / 2 - i \omega t} \) as a function of the average photon number, \( \bar{n} \), and the bandwidth, \( \Gamma \). Regions (b) and (c) are the high-energy regions, \( \bar{n} \geq 10 \). The fringes in the region (b) capture Rabi oscillations [33]. Bright fringes correspond to complete inversions of the emitter’s population, after which spontaneous emission occurs. Conversely, dark fringes correspond to complete Rabi oscillations, which leave the emitter in its ground state.

Region (a) (\( \bar{n} \leq 1 \)) is the low-energy regime on which we focus from now on. As it appears on Fig. 3(a), \( B_2^p \) never vanishes in this region, i.e. spin and light can only become partially entangled. Conversely, we consider coherent superpositions of zero and single photon states, \(|0_\uparrow\rangle + e^{i\phi} |1_\uparrow\rangle\), with \(|c_0|^2 + |c_1|^2 = 1\), \(|1_\uparrow\rangle = \int dt \xi(t)b^\dagger(t)|0\rangle\), and \( \int dt |\xi(t)|^2 = 1\). Such superpositions can be generated with high purity by semiconducting quantum dots embedded in micropillars [22].

The analytical expression of the resulting pointer states is provided in [32]. It gives rise to the QBhat further denoted \( B_2^p \), which is plotted in Fig. 3(b) for \( \xi(t) \) a decaying exponential wavepacket. Maximally entangled spin-light states can be produced for any input energy \( 1/2 \leq \bar{n} \leq 1 \), provided that \( \bar{n} = (1 + \Gamma / \gamma) / 2 \). This signals the existence of a quantum advantage, quantum superposition producing more entanglement than coherent ones with the same amount of energy.

To interpret Fig. 3(b), it is convenient to analyze H-polarized photons as a coherent superposition of circularly polarized ones \( |1_\uparrow\rangle = (|1_L\rangle + |1_R\rangle)/\sqrt{2} \). As explained above, each component of the superposition undergoes a phase shift conditioned to the spin state while being scattered. This induces the clockwise or counterclockwise rotation of the light polarization [32], the final polarizations and shapes of the pointer states depending on the energy, shape and duration of the input pulse. Therefore in general, light-matter entanglement is spread over the polarization and temporal light degrees of freedom. Two limiting cases can be considered. Mode-matched single photons (\( \Gamma \sim \gamma \)) give rise to pointer states where information is mostly encoded in the polarization of \( |\psi_1\rangle \langle \psi_1| \) that rotates from H to R (resp. L) during the interaction [32]. In this case however, the pulse shape is modified by the scattering process, altering its potential for quantum information protocols [28]. Conversely in the limit of monochromatic input wave-packets, the scattering mechanism verifies Eq. (2). On Fig. 3(b), this corresponds to \( \Gamma \ll \gamma \) and \( \bar{n} = 1/2 \). It gives rise to the pointer states \( |\psi_1\rangle = (|0\rangle + i|1\rangle)/\sqrt{2} \), with \( |1\rangle = i (|1_L\rangle - |1_R\rangle)/\sqrt{2} \), where complete information on the spin state is carried by the phase of the quantum superposition.

**Spin readout** - We now investigate if the quantum advantage appearing along entanglement generation is maintained as information on the spin is extracted at the classical level. A natural figure of merit is the distinguishability of the spin states, when the light pulses undergo a classical measurement. It is captured by the classical (cl) Bhattacharyya coefficient (cBhat) [30], which quantifies the overlap among the conditional probabilities \( p_{\uparrow|1}(x) \) of obtaining the classical outcome \( x \) among the set \( \mathcal{X} \):

\[
B_{cl} = \sum_{x \in \mathcal{X}} \sqrt{p_{\uparrow}(x)p_{\downarrow}(x)}. \tag{7}
\]

\( B_{cl} = 1 \) signals identical distributions where the measurement does not provide any information about the spin state; con-
versely $B_{cl} = 0$ corresponds to disjoint distributions granting complete knowledge of the spin state. Interestingly, the $cBhat$ is lower bounded by the $qBhat$, i.e. $B_{q} \leq B_{cl}$, with the equality being satisfied when both coefficients vanish [30].

We probe the spin with $R$-polarized, low intensity pulses, and take the Michelson interferometer depicted on Fig. 4(a) as our classical measurement apparatus. The interferometer encompases two balanced non-polarized beam splitters, each arm containing either a tunable phase-plate ($\phi$) or the SPI. A photo-counter is positioned in the output port of the interferometer. The number of detected photons $x = 0, 1$ defines our two possible classical outcomes. The phase $\phi$ is tuned, such that $p_{1}(x) = 1$ (resp. 0) if the spin is $|↑\rangle$ (resp. $|↓\rangle$) for a monochromatic single photon input pulse. Hence, the $cBhat$ $B_{cl}^{p}$ vanishes when $\Gamma/\gamma \rightarrow 0$. The overlap between the probability distributions $p_{1}(x)$ and $p_{1}(x)$ increases as soon as the phase shift induced by the SPI is reduced, or if the wave-packet shape is altered by the scattering. This is the case for short single photon pulses or coherent pulses of any duration [32].

We analyze the performance of the readout if the spin is probed with a coherent or a quantum pulse, both characterized by an exponential shape of width $\Gamma$ and a mean number of photons $\bar{n} = 1$. We introduce the light pointer states $|\Psi_{1}\rangle$ and $|\Psi_{f}\rangle$, which feature the pulses exiting the interferometer that are respectively correlated with the spin up or down. Their overlap defines the two $qBhat$ plotted on Fig. 4 as a function of the pulse bandwidth (blue lines). Their behavior is consistent with the study above, revealing a quantum advantage highlighted by the gray area and captured by $B_{q}^{qs} \geq B_{cl}^{qs}$. The difference between the $qBhat$ is maximal in the linear regime $\gamma \gg \Gamma$.

The evolution of the $cBhat$ and the $qBhat$ is plotted on Fig. 4(b) as a function of $\Gamma/\gamma$, for the two statistics of input pulses (quantum and coherent, resp. corresponding to the blue and black curves). Importantly since each $cBhat$ is lower bounded by the corresponding $qBhat$, the gray area between the blue curves is accessible to the $cBhat$ of the quantum light, but forbidden to the one of the coherent field. This signals another kind of quantum advantage, i.e. a regime where the quantum light provides a better readout than any possible classical scheme where the spin is probed with a coherent pulse of the same duration and same mean photon number.

When including pure dephasing and an efficiency reduction in the numerical model used to compute the Bhattacharyya coefficients [32], we find that the quantum advantage region is reachable for setups having an overall efficiency no less than 80%, and a dephasing rate no more than 25% of the optical decay rate $\gamma$. State-of-the-art open cavity quantum dot devices operating in the near IR can reach coupling efficiencies exceeding 90% with low dephasing rates of 0.025$\gamma$ for $\gamma^{-1}$ $\approx$ 100 ps [34], providing a promising direction towards experimental realization of the proposed setup. Commercially available superconducting nanowire detectors can also reach detection efficiencies exceeding 90% in the near IR. In principle, the required quantum input light can be produced by a similar quantum dot device. However, this would compound the degrading effects of dephasing and would also bring the overall experimental efficiency below the 80% bound. Near-term experimental realization could be possible using pulsed SPDC, with an appropriate bandwidth of at least $\Gamma = 10^{-2}\gamma$. In this case, post-selecting on successfully-created single photons, may bring the overall efficiency above the 80% bound allowing for an observation of the quantum advantage.

**Conclusions** - We studied the interaction between a spin-carrying quantum emitter and a travelling pulse of light. We considered the low-energy regime where the light carries a maximum of one excitation in average, and compared a coherent field with a quantum superposition of zero and single photon states. We find that the latter state produces spin-light entanglement more efficiently than the former, providing a quantum advantage. This quantum advantage is shown to be maintained when information on the spin...
state is extracted at the classical level, and can be observed within state-of-the-art physical implementations. Our study brings out a new interest in the exploitation of quantum resources based on energy efficiency. This inquiry is relevant from a fundamental point of view and useful to inspire new applications in the field of optical quantum computation, e.g. photon-photon gates and cluster states.

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[1] D. P. DiVincenzo and IBM, The Physical Implementation of Quantum Computation, arXiv:quant-ph/0002077, 10.1022/1521-3978/200009/48/9/11/771::AID-PROP771>3.0.CO;2-E (2000), arXiv: quant-ph/0002077.
[2] H. J. Kimble, The quantum internet, Nature 453, 1023–1030 (2008).
[3] H-J. Briegel, W. Dür, J. I. Cirac, and P. Zoller, Quantum Repeaters: The Role of Perfect Local Operations in Quantum Communication, Physical Review Letters 81, 5932 (1998).
[4] T. E. Northup and R. Blatt, Quantum information transfer using photons, Nature Photonics 8, 356 (2014).
[5] F. Flamini, N. Spagnolo, and F. Sciarrino, Photonic quantum information processing: a review, Reports on Progress in Physics 82, 016001 (2019).
[6] W. B. Gao, P. Fallah, E. Togan, J. Miguel-Sanchez, and A. Imaeorglu, Observation of entanglement between a quantum dot spin and a single photon, Nature 491, 426 (2012).
[7] K. De Greve, P. L. McMahon, L. Yü, J. S. Pelc, C. Jones, C. M. Natarajan, N. Y. Kim, E. Aebi, S. Maier, C. Schneider, M. Kamp, S. Hofling, R. H. Hadfield, A. Forchel, M. M. Fejer, and Y. Yamamoto, Complete tomography of a high-fidelity solid-state spin photon qubit pair, Nature Communications 4, 2228 (2013).
[8] J. R. Schaibley, A. P. Burgers, G. A. McCracken, L.-M. Duan, P. R. Berman, D. G. Steel, A. S. Bracker, D. Gammon, and L. J. Sham, Demonstration of Quantum Entanglement between a Single Electron Spin Confined to an InAs Quantum Dot and a Photon, Physical Review Letters 110, 167401 (2013).
[9] J. Lee, B. Villa, A. Bennett, R. Stevenson, D. Ellis, I. Farrer, D. Ritchie, and A. Shields, A Quantum dot as a source of time-bin entangled multi-photon states, Quantum Science and Technology 4, 025011 (2019).
[10] A. Tchebotareva, S. L. N. Hermans, P. C. Humphreys, D. Voigt, P. J. Harnsma, L. K. Cheng, A. L. Verlaan, N. Dijkhuizen, W. de Jong, A. Drué, and R. Hanson, Entanglement between a diamond spin qubit and a photonic time-bin qubit at telecom wavelength, Phys. Rev. Lett. 123, 063601 (2019).
[11] A. Javadi, D. Ding, M. H. Appel, S. Mahmoodian, M. C. Löbl, I. Söllner, R. Schott, C. Papon, T. Pregnolato, S. Stobbe, L. Middo, T. Schröder, A. D. Wieck, A. Ludwig, R. J. Warburton, and P. Lodahl, Spin–photon interface and spin-controlled photon switching in a nanobeam waveguide, Nature Nanotechnology 13, 398 (2018).
[12] N. H. Lindner and T. Rudolph, Proposal for Pulsed On-Demand Sources of Photonic Cluster State Strings, Physical Review Letters 103, 113602 (2009).
[13] I. Schwartz, D. Cogan, E. R. Schmidgall, Y. Don, L. Gantz, O. Kenneth, N. H. Lindner, and D. Gerbshoni, Deterministic Generation of a Cluster State of Entangled Photons, Science 354, 434 (2016), arXiv: 1606.07492.
[14] C. Y. Hu, W. J. Munro, and J. G. Rarity, Deterministic photon entangler using a charged quantum dot inside a microcavity, Physical Review B (Condensed Matter and Materials Physics) 78, 125318 (2008).
[15] C. Y. Hu, A. Young, J. L. O’ Brien, W. J. Munro, and J. G. Rarity, Giant optical faraday rotation induced by a single-electron spin in a quantum dot: Applications to entangling remote spins via a single photon, Phys. Rev. B 78, 085307 (2008).
[16] M. N. Leuenberger, M. E. Flatté, and D. D. Awschalom, Teleportation of electronic many-qubit states encoded in the electron spin of quantum dots via single photons, Physical Review Letters 94, 107401 (2005).
[17] C. Bonato, F. Haapt, S. S. R. Oemrawsingh, J. Gudat, D. Ding, M. P. van Exter, and D. Bouwmeester, CNOT and Bell-state analysis in the weak-coupling cavity QED regime, Physical Review Letters 104, 160503 (2010).
[18] V. Giesel, N. Somaschi, G. Hornecker, T. Grange, B. Reznchenko, L. D. Santis, J. Demory, C. Gomez, I. Sagnes, A. Lemaitre, O. Krebs, N. D. Lanzillotti-Kimura, L. Lanco, A. Affeves, and P. Senellart, Coherent manipulation of a solid-state artificial atom with few photons, Nature Communications 7, 10.1038/ncomms11986 (2016).
[19] L. D. Santis, C. Antón, B. Reznchenko, N. Somaschi, G. Coppola, J. Senellart, C. Gómez, A. Lemaitre, I. Sagnes, A. G. White, L. Lanco, A. Affeves, and P. Senellart, A solid-state single-photon filter, Nature Nanotechnology 12, 663 (2017).
[20] M. Maffei, P. A. Camati, and A. Affeves, Closed-System Solution of the 1D Atom from Collision Model, arXiv:2112.09672 [quant-ph] (2021), arXiv: 2112.09672.
[21] M. Maffei, P. A. Camati, and A. Affeves, Probing nonclassical light fields with energetic witnesses in waveguide quantum electrodynamics, Physical Review Research 3, L032073 (2021).
[22] J. Loredo, C. Antón, B. Reznchenko, P. Hilaire, A. Harouri, C. Millet, H. Ollivier, N. Somaschi, L. De Santis, A. Lemaitre, et al., Generation of non-classical light in a photon-number superposition, Nature Photonics 13, 803 (2019).
[23] C. Y. Hu, Spin-based single-photon transistor, dynamic random access memory, diodes, and routers in semiconductors, Phys. Rev. B 94, 245307 (2016), arXiv:1704.02610 [cond-mat.mes-hall].
[24] J. A. Gross, C. M. Caves, G. J. Milburn, and J. Combes, Qubit models of weak continuous measurements: Markovian conditional and open-system dynamics, Quantum Science and Technology 3, 024005 (2018).
[25] R. Loudon, The Quantum Theory of Light (OUP Oxford, 2000).
[26] C. W. Gardiner and M. J. Collett, Input and output in damped quantum systems: Quantum stochastic differential equations and the master equation, Physical Review A 31, 3761 (1985).
[27] F. Ciccarello, Collision models in quantum optics, Quantum Science and Technology 3, 024011 (2018).
[28] K. Kojima, H. F. Hofmann, S. Takeuchi, and K. Sasaki, Efficiencies for the single-mode operation of a quantum optical nonlinear shift gate, Phys. Rev. A 70, 013810 (2004).
[29] W. H. Zurek, Decoherence, einselection, and the quantum origins of the classical, Reviews of Modern Physics 75, 715 (2003), publisher: American Physical Society.
[30] C. A. Fuchs and J. van de Graaf, Cryptographic Distinguishability Measures for Quantum Mechanical States, arXiv:quant-ph/9712042 (1998), arXiv: quant-ph/9712042.
[31] C. A. Fuchs and C. M. Caves, Mathematical Techniques for Quantum Communication Theory, arXiv:quant-ph/9604001 (1996), arXiv: quant-ph/9604001.
[32] See Supplemental Material at [URL will be inserted by publisher] for further details.
[33] M. O. Scully and M. S. Zubairy, Quantum Optics (1997).
[34] N. Tomm, A. Javadi, N. O. Antoniadis, D. Najer, M. C. Löbl, A. R. Korsch, R. Schott, S. R. Valentin, A. D. Wieck, A. Ludwig, and et al., A bright and fast source of coherent single photons, Nature Nanotechnology 16, 399 (2021).