Synthesis of arbitrary Fock states via conditional measurement on beam splitters

B. M. Escher,[1] A. T. Avelar,[1,2,*] and B. Baseia[1]
[1]*Instituto de Física, Universidade Federal de Goiás, 74.001-970, Goiânia (GO), Brazil*
[2]*Instituto de Física, Universidade de Brasília, 70.919-970, Brasília (DF), Brazil*
(Received 3 September 2004; revised manuscript received 28 June 2005; published 13 October 2005)

In a previous work [Opt. Commun. **138**, 71 (1997)] a scheme was proposed to create traveling fields in the Fock state $|2^J\rangle$. Here we show how to extend this result to arbitrary Fock states. The procedure combines one-photon states impinging on a sequence of distinct beam splitters, each one associated with a (zero detection) single-photon photodetector, with optimization of the success probability to get the desired state. Advantages and disadvantages of this scheme are discussed.

DOI: 10.1103/PhysRevA.72.045803 PACS number(s): 42.50.Dv

Many types of nonclassical states have been obtained in laboratories in the last years; representative examples of them being the squeezed state [1], the "Schrodinger's cat" state [2], the displaced single-photon Fock state [3], and the state $c_0|0\rangle+c_1|1\rangle$ [4]. Various proposals to create arbitrary and specific states, for trapped fields inside high-Q cavities [5–8] and traveling fields [9–16], appeared in the literature. To give some examples of specific states required for measuring properties of other fields, we mention: (i) the *reciprocal binomial state* [17], required for the measurement of the phase distribution $P(\theta)$ of an arbitrary state [18], and for quantum lithography [19]; (ii) the *complementary coherent state* [20], to obtain the Husimi Q-function of arbitrary states [21]; (iii) the *polynomial state* [22,23] to measure variances of quadrature operators [21], all of them using the projection synthesis scheme [18]; (iv) the *vacuum squeezed state* to measure the phase distribution $P(\theta)$ using the multiport scheme, as investigated in [14,24]; (v) the *coherent state*, to measure the coefficients of arbitrary states using the balanced multiport technique [25]; etc.

Another important state studied in the literature is the Fock state, a kind of "work horse" of quantum optics. It has relevant potential applications on secure quantum communication [26], quantum cryptography [27], optimal capacity coding in quantum channels [28], in high-precision quantum interferometry [29], and in reconstruction of density operators [30]. However, its generation in laboratories is not trivial, particularly for highly excited Fock states. This has motivated proposals in this direction, for stationary [6] and traveling fields [10,15,16]. In the later case an interesting proposal is one by Steuernagel [10]. The scheme uses an array of beam splitters with several input ports which are fed with single-photon Fock states by conditional measurement in all output ports, except for the last one, from which the desired state emerges. The method assumes all beam splitters (BS) with 50-50 transmission coefficients to generate the family of Fock states $|2^J\rangle$, $J=1, 2, 3, \ldots$ with maximum success probability. While the scheme in [10] employs only 50-50 beam splitters, in our scheme below all BS have distinct transmission coefficients. So, in [10] the same BS can be recycled whereas in our scheme this is not possible. In this regard [10] is much nicer than the idea presented here. However, this later is better than that in [10] in respect to the applicability: while our idea is applicable to arbitrary number states, Ref. [10] only allows one to synthesize the number states $|2^J\rangle$.

Our alternative scheme is inspired by that of Dakna *et al.* [11] and is reminiscent of another one preparing a photon-added state [31]. It generalizes the result found in Ref. [10] for arbitrary Fock states $|N\rangle$, N integer. To this end Fig. 1 stands for the first step of the entire arrangement [32].

In this figure the role played by the beam splitter (BS) upon the input state $|\Psi\rangle_{in}=|1\rangle_a|1\rangle_{b_1}$ results in the output state $|\Psi\rangle_{\text{out\_1}}=\hat{R}_1|\Psi\rangle_{\text{in}}$, where $\hat{R}_1$ corresponds to the (unitary) operator [33]

$$\hat{R}_1 = \exp[i\theta_1(\hat{a}^\dagger\hat{b}_1 + \hat{a}\hat{b}_1^\dagger)], \tag{1}$$

in terms of the creation (annihilation) operators $\hat{a}^\dagger(\hat{a})$ and $\hat{b}_1^\dagger(\hat{b}_1)$ for modes a and b_1. Using Eq. (1) one obtains the following transformations:

$$\hat{R}_1 a^\dagger \hat{R}_1^\dagger = t_1\hat{a}^\dagger + r_1\hat{b}_1^\dagger, \quad \hat{R}_1 b_1^\dagger \hat{R}_1^\dagger = r_1\hat{a}^\dagger + t_1\hat{b}_1^\dagger, \tag{2}$$

and the application of the operator $\hat{R}_1$ upon the input state $|\Psi\rangle_{\text{in}}$ leads to the output state

$$\begin{aligned} |\Psi\rangle_{\text{out\_1}} &= \hat{R}_1|1\rangle_a|1\rangle_{b_1} = \hat{R}_1\hat{a}^\dagger\hat{R}_1^\dagger\hat{R}_1\hat{b}^\dagger\hat{R}_1^\dagger|0\rangle_a|0\rangle_{b_1} \\ &= (t_1\hat{a}^\dagger + r_1\hat{b}_1^\dagger)(r_1\hat{a}^\dagger + t_1\hat{b}_1^\dagger)|0\rangle_a|0\rangle_{b_1}, \end{aligned} \tag{3}$$

where $t_1=\cos(\theta_1)$ and $r_1=i\sin(\theta_1)$ are the transmittance and reflectance of the BS_1, respectively. A Fock state is obtained in the a-mode when an appropriate conditional measurement

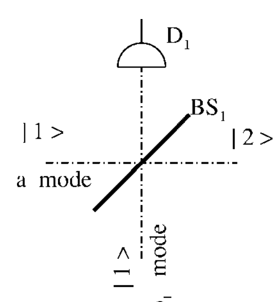

FIG. 1. Schematical setup generating the Fock state $|2\rangle$.

*Corresponding author. Email address: ardiley@unb.br

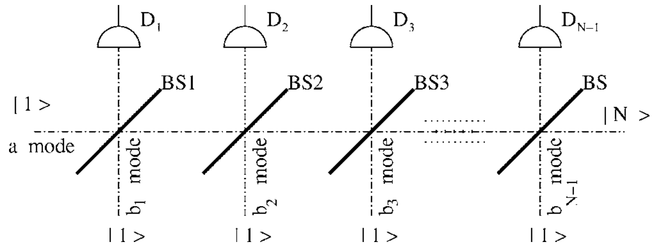

FIG. 2. Iterative procedure creating the Fock state $|N\rangle$ using single Fock states in the a- and b_j-modes, $j=1, 2, 3, \ldots$

is made, namely, when no photon is detected in the output b_1-mode, as follows,

$$|\phi^{(2)}\rangle = {}_b\langle 0|\Psi\rangle_{\text{out\_1}} = r_1 t_1 (\hat{a}^\dagger)^2 |0\rangle_a, \tag{4}$$

which corresponds to the (unnormalized) Fock state $|2\rangle$.

The same procedure is valid for the next step, when the state $|2\rangle$ arrives at the BS2 in the a-mode with the state $|1\rangle$ in the b_2 -mode, with no photon detected in the output b_2-mode, resulting in

$$|\phi^{(3)}\rangle = {}_{b_2}\langle 0|\Psi\rangle_{\text{out\_2}} = \frac{r_2 t_2^2}{\sqrt{2}} (\hat{a}^\dagger)^3 |0\rangle_a. \tag{5}$$

Repeating this procedure $N-1$ times we obtain the arbitrary Fock state $|N\rangle$

$$|\phi^{(N)}\rangle = {}_{b_{N-1}}\langle 0| \sum_{n=0}^{N-1} \frac{\binom{N-1}{n}}{\sqrt{N!}} t_k^{N-n-1} r_k^n (\hat{a}^\dagger)^{N-n-1} (\hat{b}_k^\dagger)^n$$
$$\times (r_k \hat{a}^\dagger + t_k \hat{b}_k^\dagger) |0\rangle_a |0\rangle_{b_{N-1}} \tag{6}$$

where $t_k = \cos(\theta_k)$ and $r_k = i\sin(\theta_k)$ are the transmittance and reflectance of the kth BS, respectively. The result in Eq. (6) corresponds to the arrangement displayed in Fig. 2.

Now, the success probability to get the Fock state $|2\rangle$ from the state $|1\rangle$ is given by $P_{|2\rangle} = \langle \phi^{(2)} | \phi^{(2)} \rangle$, resulting in $P_{|2\rangle} = 2(1-t_1^2)t_1^2$, which attains the maximum value $P_{|2\rangle}^{(\max)} = 1/2$ for $t_1 = \sqrt{1/2}$. The same calculation for the BS2 leads to $P_{|3\rangle}^{(\max)} = 4/9$, with $t_2 = \sqrt{2/3}$. So, for the kth step the success probability reads

$$P_{|k+1\rangle} = (k+1)(1-t_k^2) t_k^{2k}, \tag{7}$$

which attains the maximum value

$$P_{|k+1\rangle}^{(\max)} = \left(\frac{k}{k+1} \right)^k, \tag{8}$$

when we choose the transmission coefficient t_k as

$$t_k = \left(\frac{k}{k+1} \right)^{1/2}. \tag{9}$$

Hence, the total success probability $\mathcal{P}$ coming from the product of those in all steps shown in Fig. 2 results in

$$\mathcal{P}_{|N\rangle}^{(\max)} = \prod_{k=1}^{N-1} P_{|k+1\rangle}^{(\max)} = \prod_{k=1}^{N-1} \left[\left(\frac{k}{k+1} \right)^k \right] = \frac{N!}{N^N}, \tag{10}$$

which coincides with the maximum success probability found in Ref. [10]

The present approach is a convenient adaptation of a previous one by Dakna *et al.* [11]. A comparison between them shows that they have some analogies and some differences: they are similar in the use of the same (array) configuration of BS, with conditional measurements; they differ in the use of distinct states impinging the a-mode in Fig. 2 (Ref. [11] employs the vacuum state as the initial state whereas we have employed a single-photon Fock state, which economizes one detector) and the use of BS with distinct transmittances for the optimization of success probability to get the state (in Ref. [11] the same transmittance for all BS is assumed). In both cases ideal photodetectors (unity efficiency) has been employed; recent technological advances have achieved photodetectors with efficiency near 100% [34] improving the fidelity of the state being created. In addition, while in [11] it was verified that maximum success probability to create the truncated coherent phase state is not substantially affected when using either equal or different BS, here we observe that this result is not true for the creation of Fock states: for example, the maximum success probability to obtain the state $|2\rangle$ results is 67% (20%) when using distinct (identical) BS. This result remains valid for $N=3, 4, 5, \ldots$.

Until now, we have assumed indistinguishable photons on demand impinging on the input of all BS. The "observability of two-photon interference effects naturally requires that two single photons arriving at the two input ports of the beam splitter be indistinguishable in terms of their pulsewidth, bandwidth, polarization, carrier frequency, and arrival time (time-jitter) at the beam splitter" [35]. This implies that the spatial-temporal mode profiles of all the single photons that are supposed to be added up have to be "identical"; otherwise interference effects occurring at each BS are reduced or even destroyed [36]. For example, if two *distinguishable* photons impinge the BS they can behave independently and the two-photon interference effect is reduced [32,37]. The identity of wave packets also constitute a crucial requirement in quantum information processing [38] and in quantum networking [39].

Nowdays, there exist stable sources on demand yielding many single photons on the input ports of the many beam splitters, e.g., using single quantum dots [40]; single molecules [41]; diamond colour centers [42]; atoms [43]; turnstile device [44]; and parametric down conversion schemes [45]. However, to date, such sources are not able to attain completely the mentioned photon properties. For example, in single atom or quantum dot sources the pulsewidth and bandwidth will not distinguish the interfering photons, but the uncertainity in photon emission (*time-jitter*) will. Also, while various sources of distinguishability can be eliminated, the inherent jitter in photon emission time remains as an unavoidable source of distinguishability [35]. The time delay (time-jitter) associates a temporal phase to the interfering photons (not a phase of the single photon state, which has no definite phase [32]). This would constitute a problem affecting all schemes (e.g., [8,10,11]) using many single photon sources to create states of traveling fields.

To verify the robustness of our scheme against deleterious effect of the time jitter, we analyze the reduction of fidelity of the state created. In the presence of a temporal jitter ϵ, in the photon emission time a single photon state in the b-mode is represented as [35]

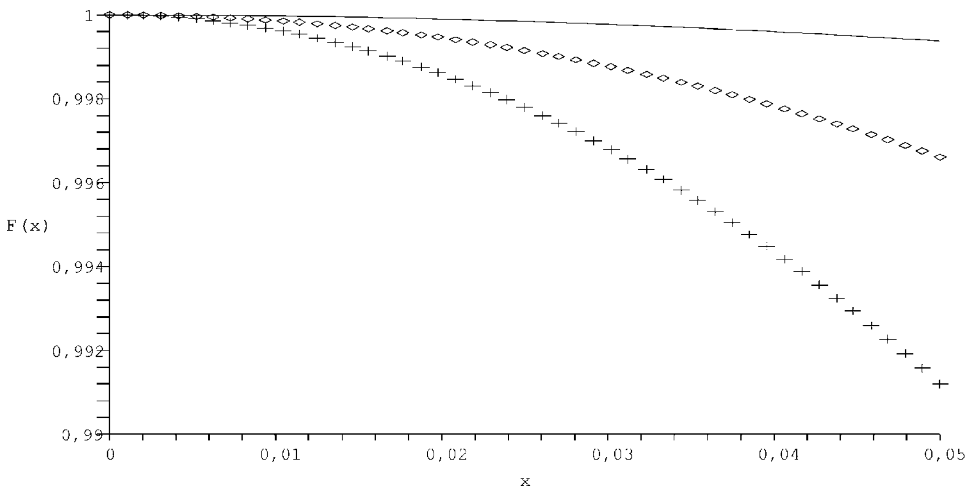

FIG. 3. Plots of the fidelity $F(x)$ for the generation of the states $|2\rangle$ (solid), $|5\rangle$ (dotted), and $|10\rangle$ (cross), as a function of time jitter normalized to photon pulse width $x=\epsilon/\tau$.

$$|1\rangle_b = \int d\omega f(\omega) e^{i\omega\epsilon}\hat{b}^\dagger|0\rangle_b, \qquad (11)$$

where $f(\omega)=(\tau^2/\pi)^{1/4}e^{-(\omega-\omega_0)^2\tau^2/2}$ is the spectrum of the photon wave packet, with $\int f^2 d\omega=1$. For clarity, we keep all photons in the remaining modes as ideal and indistiguishable. So, the initial state $|\psi\rangle_{in}=|N\rangle_a|1\rangle_b$ in the last BS will be written as

$$|\psi\rangle_{\text{in}}^* = \frac{1}{\sqrt{N!}}\int d\omega_1\cdots d\omega_N f(\omega_1)\cdots f(\omega_N)\hat{a}^\dagger(\omega_1)\cdots\hat{a}^\dagger(\omega_N)$$

$$\times\int d\omega_{N+1}f(\omega_{N+1})e^{i\epsilon\omega_{N+1}}\hat{b}^\dagger(\omega_{N+1})|0\rangle_a|0\rangle_b. \qquad (12)$$

Using Eq. (2), the commutation relation $[\hat{a}(\omega_i),\hat{a}^\dagger(\omega_j)] =\delta(\omega_i-\omega_j)$ and assuming no detection events in the b-mode, the output state in a-mode $|\phi^{(N+1)}\rangle^* = {}_b\langle 0|\psi\rangle_{in}^*$ collapses to

$$|\phi^{(N+1)}\rangle^* = \frac{RT^N}{\sqrt{N!}}\int d\omega_1\cdots d\omega_{N+1}f(\omega_1)\cdots f(\omega_{N+1})$$

$$\times e^{i\epsilon\omega_{N+1}}\hat{a}^\dagger(\omega_1)\cdots\hat{a}^\dagger(\omega_{N+1})|0\rangle. \qquad (13)$$

Setting the time jitter $\epsilon=0$ the success probability $\mathcal{P} = {}^*\langle\phi^{(N+1)}|\phi^{(N+1)}\rangle^*$ coincides with the result given by Eq. (7). Equation (13) allows us to find the effect of time jitter upon the fidelity. We obtain,

$$F(\epsilon)=\prod_{k=1}^{N}\frac{|\langle\phi^{(k)}|\phi^{(k)}\rangle^*|^2}{\||\phi^{(k)}\rangle\|^2\||\phi^{(k)}\rangle^*\|^2} = \frac{(N+1)!}{\prod_{k=1}^{N}(ke^{(\epsilon/\tau)^2/2}+1)}. \qquad (14)$$

Figure 3 shows the plots of fidelity versus normalized time-jitter for the Fock states $|2\rangle$, $|5\rangle$, $|10\rangle$. Note that fidelity 100% results for $\epsilon=0$.

For an incoherently pumped quantum dot single photon source the emission time-jitter is on the order of 1×10^{-11} s and for a pulse width about 1×10^{-9} s results the limiting rate $\epsilon/\tau\cong 0.01$ [35] which yields for our case fidelity greater than 99%, as is evident in Fig. 3.

In conclusion, we have extended the results obtained in Ref. [10], creating Fock states of kind $|2^J\rangle$, $J=1, 2, 3, \ldots$, to arbitrary Fock states $|N\rangle$, $N=1, 2, 3, \ldots$. In both cases the maxima success probabilities to get a desired state coincide. To obtain this extension we have employed a set of BS in distinct configurations, with the BS having different transmission coefficients t_k[cf. Eq. (9)]; they are all equal (50-50) in Ref. [10]. Here it is worth stressing that all values t_k required to generate a state $|N\rangle$ with maximum success probability do not change when one passes to the next step $|N+1\rangle$, except for the additional Nth BS. It is worth stressing, for comparison, that in a recent Rapid Communication [16] the Fock state $|5\rangle$ was proposed via linear optical extraction from coherent state with a success probability of 10^{-10}% (fidelity of 79%), whereas in our scheme it results are 4% (100%). Finally, we have considered the effect of time jitter upon the fidelity of our states being created. This (deleterious) effect depends on the generation scheme, on the type of photon source employed, and the wanted state. It was shown that, in our scheme a fidelity greater than 99% is achieved using quantum dot sources to create Fock states (Fig. 3).

We thank Dr. C. H. Monken for helpful references; the support from the CNPq and FUNAPE (GO), Brazilian agencies, are also acknowledged.

[1] R. E. Slusher, L. W. Hollberg, B. Yurke, J. C. Mertz, and J. F. Valley, Phys. Rev. Lett. **55**, 2409 (1985).

[2] M. Brune, E. Hagley, J. Dreyer, X. Maitre, A. Maali, C. Wunderlich, J. M. Raimond, and S. Haroche, Phys. Rev. Lett. **77**, 4887 (1996).

[3] A. I. Lvovsky and S. A. Babichev, Phys. Rev. A **66**, 011801(R) (2002).

[4] A. I. Lvovsky and J. Mlynek, Phys. Rev. Lett. **88**, 250401 (2002); S. A. Babichev, J. Ries, and A. I. Lvovsky, Europhys. Lett. **64**, 1 (2003).

[5] K. Vogel, V. M. Akulin, and W. P. Schleich, Phys. Rev. Lett. **71**, 1816 (1993); A. S. Parkins, P. Marte, P. Zoller, and H. J. Kimble, *ibid.* **71**, 3095 (1993); B. M. Garraway, B. Sherman, H. Moya-Cessa, P. L. Knight, and G. Kurizki, Phys. Rev. A **49**, 535 (1994); C. K. Law and J. H. Eberly, Phys. Rev. Lett. **76**, 1055 (1996).

[6] P. Domokos, J. Janszky, and P. Adam, Phys. Rev. A **50**, 3340 (1994); B. T. H. Varcoe, S. Brattke, M. Weidinger, and H. Walther, Nature (London) **403**, 743 (2000); S. Brattke, B. T. H. Varcoe, and H. Walther, Phys. Rev. Lett. **86**, 3534 (2001); M. F. Santos, E. Solano, and R. L. de Matos Filho, *ibid.* **87**, 093601 (2001).
[7] L. A. de Souza, B. Baseia, and J. M. C. Malbouisson, Phys. Lett. A **309**, 5 (2003); A. T. Avelar, L. A. de Souza, T. M. Rocha Filho, and B. Baseia, J. Opt. B: Quantum Semiclassical Opt. **6**, 383 (2004); A. Aragão, A. T. Avelar, J. M. C. Malbouisson, and B. Baseia, Phys. Lett. A **329**, 284 (2004).
[8] L. P. A. Maia, A. T. Avelar, J. M. C. Malbouisson, and B. Baseia, J. Opt. B: Quantum Semiclassical Opt. **6**, 351 (2004); P. B. Monteiro, B. Baseia, A. T. Avelar, and J. M. C. Malbouisson, Braz. J. Phys. **35**, 579 (2005).
[9] D. T. Pegg, L. S. Phillips, and S. M. Barnett, Phys. Rev. Lett. **81**, 1604 (1998); M. G. A. Paris, Phys. Rev. A **62**, 033813 (2000).
[10] O. Steuernagel, Opt. Commun. **138**, 71 (1997).
[11] M. Dakna, J. Clausen, L. Knoll, and D. G. Welsch, Phys. Rev. A **59**, 1658 (1999); **60**, 726 (1999).
[12] C. J. Villas-Boas, Y. Guimarães, M. H. Y. Moussa, and B. Baseia, Phys. Rev. A **63**, 055801 (2001).
[13] A. T. Avelar, B. M. Escher, and B. Baseia, Phys. Lett. A **318**, 161 (2003); A. T. Avelar, T. M. Rocha Filho, and B. Baseia, Opt. Commun. **231**, 297 (2004); A. T. Avelar, J. M. C. Malbouisson, and B. Baseia, Physica A **334**, 139 (2004); B. M. Escher, A. T. Avelar, T. M. da Rocha Filho, and B. Baseia, Phys. Rev. A **70**, 025801 (2004).
[14] X. Zou, K. Pahlke, and W. Mathis, Phys. Lett. A **323**, 329 (2004).
[15] S. Ya. Kilin and D. B. Horoshko, Phys. Rev. Lett. **74**, 5206 (1995); G. M. D'Ariano, L. Maccone, M. G. A. Paris, and M. F. Sacchi, Phys. Rev. A **61**, 053817 (2000).
[16] K. Sanaka, Phys. Rev. A **71**, 021801(R) (2005).
[17] C. Valverde, A. T. Avelar, B. Baseia, and J. M. C. Malbouisson, Phys. Lett. A **315**, 213 (2003).
[18] S. M. Barnett and D. T. Pegg, Phys. Rev. Lett. **76**, 4148 (1996).
[19] See,e.g., G. Björk, L. L. Sanchez-Soto, and J. Soderholm, Phys. Rev. Lett. **86**, 4516 (2001).
[20] A. T. Avelar, B. Baseia, and N. G. de Almeida, J. Opt. B: Quantum Semiclassical Opt. **6**, 41 (2004).
[21] B. Baseia, M. H. Y. Moussa, and V. S. Bagnato, Phys. Lett. A **231**, 4148 (1997).
[22] J. A. Bergou, M. Hillery, and D. Yu, Phys. Rev. A **43**, 515 (1991).
[23] S. Souza, A. T. Avelar, N. G. de Almeida, and B. Baseia, Opt. Commun. **239**, 281 (2004).
[24] K. L. Pregnell and D. T. Pegg, Phys. Rev. Lett. **89**, 173601 (2002).
[25] H. Paul, P. Torma, T. Kiss, and I. Jex, Phys. Rev. Lett. **76**, 2464 (1996).
[26] S. J. van Enk, J. I. Cirac, and P. Zoller, Phys. Rev. Lett. **78**, 4293 (1997); K. M. Gheri, C. Saavedra, P. Torma, J. I. Cirac, and P. Zoller, Phys. Rev. A **58**, R2627 (1998); S. J. van Enk, J. I. Cirac, and P. Zoller, Science **279**, 205 (1998).
[27] T. Jennewein, C. Simon, G. Weihs, H. Weinfurter, and A. Zeilinger, Phys. Rev. Lett. **84**, 4729 (2000).
[28] C. M. Caves and P. D. Drummond, Rev. Mod. Phys. **66**, 481 (1994).
[29] M. J. Holland and K. Burnett, Phys. Rev. Lett. **71**, 1355 (1993).
[30] O. Steuernagel and J. A. Vaccaro, Phys. Rev. Lett. **75**, 3201 (1995).
[31] M. Dakna, L. Knöll, and D.-G. Welsch, Opt. Commun. **145**, 309 (1998).
[32] C. K. Hong, Z. Y. Ou, and L. Mandel, Phys. Rev. Lett. **59**, 2044 (1987).
[33] B. Yurke, S. L. McCall, and J. R. Klauder, Phys. Rev. A **33**, 4033 (1985); S. Prasad, M. O. Scully, and W. Martienssen, Opt. Commun. **62**, 139 (1987); R. A. Campos, B. E. A. Saleh, and M. C. Teich, Phys. Rev. A **40**, 1371 (1989).
[34] A. Imamoglu, Phys. Rev. Lett. **89**, 163602 (2002); D. F. V. James and P. G. Kwiat, *ibid.* **89**, 183601 (2002).
[35] A. Kiraz, M. Atatüre, and A. Immamoglu, Phys. Rev. A **69**, 032305 (2004).
[36] A. Imamoglu, Physica E (Amsterdam) **16**, 47 (2003).
[37] E. Moreau, I. Robert, J. M. Gérard, I. Abram, L. Manin, and V. Thierry-Mieg, Appl. Phys. Lett. **79**, 2865 (2001); M. Pelton, C. Santori, J. Vuckovic, B. Y. Zhang, G. S. Solomon, J. Plant, and Y. Yamamoto, Phys. Rev. Lett. **89**, 233602 (2002); C. Santori, D. Fattal, J. Vuckovic, G. S. Solomon, and Y. Yamamoto, New J. Phys. **6**, 89 (2004).
[38] E. Knill, R. Laflamme, and G. J. Milburn, Nature (London) **409**, 46 (2001).
[39] C. Monroe, Nature (London) **416**, 238 (2002).
[40] P. Michler, A. Kiraz, C. Becher, W. V. Schoenfeld, P. M. Petroff, L. D. Zhang, E. Hu, and A. Imamoglu, Science **290**, 2282 (2000); C. Santori, M. Pelton, G. Solomon, Y. Dale, and Y. Yamamoto, Phys. Rev. Lett. **86**, 1502 (2001); V. Zwiller, H. Blom, P. Jonsson, N. Panev, S. Jeppesen, T. Tsegaye, E. Goobar, M. E. Pistol, L. Samuelson, and G. Bjork, Appl. Phys. Lett. **78** 2476 (2001); C. Santori, D. Fattal, J. Vuckovic, G. S. Solomon, and Y. Yamamoto, Nature (London) **419**, 594 (2002); Z. Yuan, B. E. Kardynal, R. M. Stevenson, A. J. Shields, C. J. Lobo, K. Cooper, N. S. Beattie, D. A. Ritchie, and M. Pepper, Science **295**, 102 (2002).
[41] B. Lounis and W. E. Moerner, Nature (London) **407**, 491 (2000); C. Brunel, B. Lounis, P. Tamarat, and M. Orrit, Phys. Rev. Lett. **83**, 2722 (1999); F. DeMartini, G. DiGiuseppe, and M. Marrocco, *ibid.* **76**, 900 (1996).
[42] C. Kurtsiefer, S. Mayer, P. Zarda, and H. Weinfurter, Phys. Rev. Lett. **85**, 290 (2000); R. Brouri, A. Beveratos, J. P. Poizat, and P. Grangier, Opt. Lett. **25**, 1294 (2000); A. Beveratos, R. Brouri, T. Gacoin, J. P. Poizat, and P. Grangier, Phys. Rev. A **64**, 061802(R) (2001); A. Beveratos, S. Kuhn, R. Brouri, T. Gacoin, J. P. Poizat, and P. Grangier, Eur. Phys. J. D **18**, 191 (2002).
[43] A. Kuhn, M. Hennrich, and G. Rempe, Phys. Rev. Lett. **89** , 067901 (2002).
[44] J. Kim, O. Benson, H. Kan, and Y. Yamamoto, Nature (London) **397**, 500 (1999).
[45] C. K. Hong and L. Mandel, Phys. Rev. Lett. **56**, 58 (1986).