Optimal cloning of single-photon polarization by coherent feedback of beam splitter losses

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Abstract. Light fields can be amplified by measuring the field amplitude reflected at a beam splitter of reflectivity $R$ and adding a coherent amplitude proportional to the measurement result to the transmitted field. By applying the quantum optical realization of this amplification scheme to single-photon inputs, it is possible to clone the polarization states of photons. We show that optimal cloning of single-photon polarization is possible when the gain factor of the amplification is equal to $1/\sqrt{1 - R}$.

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1. Introduction: quantum cloning and field amplification

One of the fundamental features of quantum information is that it is impossible to generate perfect copies (or ‘clones’) of an unknown quantum state input \[1\]. This no-cloning theorem is particularly interesting in the light of the wave–particle dualism of optics, since the amplitude of a classical wave can be copied perfectly by any classical amplification process. Soon after the first formulation of the no-cloning theorem, it was pointed out that perfect cloning by phase sensitive optical amplification is prevented by the unavoidable spontaneous emission in such processes \[2, 3\]. However, it was later found that stimulated emission is in fact an optimal approximation to perfect quantum cloning \[4\]. This insight was quickly followed by the first experimental realizations of optical quantum cloning using parametric optical amplification \[5–8\]. Recently, it has also been discovered that the bunching properties of light fields can be used to obtain optimal clones by post-selecting the output of a beam splitter \[9\]. In general, optical cloning methods thus exploit the natural wave–particle dualism of light to clone the quantum coherence of photons by manipulating the (classical) optical coherence of the light field.

A more direct way to access the field properties of photons is to measure the quadrature components \(\hat{x}\) and \(\hat{y}\) of the complex field amplitude, \(\hat{a} = \hat{x} + i\hat{y}\). As demonstrated by a number of experimental results \[10–14\], such measurements provide quantum mechanically precise information on the coherent field properties associated with photon number states. It seems obvious that this method can also be used to measure the polarization state of a photon, since the polarization of light is completely described by the two complex amplitudes \(\hat{a}_H\) and \(\hat{a}_V\) of a pair of orthogonal polarizations H and V. For a single-photon input, the measurement of the two complex amplitudes \(\hat{a}_H\) and \(\hat{a}_V\) by homodyne detection is indeed equivalent to a quantum mechanically precise detection of the photon in the polarization defined by the measurement results obtained for the amplitudes. A particularly simple cloning scheme could thus be realized by measuring the complex amplitudes of the input photon and modulating a coherent laser beam to emit multiple photons with the same polarization amplitudes.

However, homodyne detection can also be applied to fields of unknown photon number. It is then possible to obtain partial information about the polarization of a photon by ‘dividing’ the one-photon input at a beam splitter of reflectivity \(R\) and measuring only the reflected fraction of the light. The resulting losses caused by the reduction of the transmitted amplitude by a factor of \(\sqrt{1 - R}\) can be compensated by adding a coherent laser amplitude proportional to the measurement results for the coherent amplitudes \(\hat{a}_H\) and \(\hat{a}_V\) \[15\]. It is also possible to over-compensate the losses to achieve an amplification of the field variables. In fact, it has been shown that this kind of over-compensation can be used to achieve the noiseless amplification of a single quadrature component of the light field \[16\], and the application of this scheme to the continuous variable cloning of Gaussian states has recently been demonstrated \[17, 18\]. It is thus clear that a minimal noise amplification of the light field can be achieved by a finite resolution measurement of the amplitude and an appropriate coherent feedback.

In the following, it is shown that the kind of optical amplification used to clone continuous variable field states of a single mode in \[17\] can also be used to clone the polarization state of a single-photon input. By using an optimized gain factor of \(1/\sqrt{1 - R}\) to minimize the noise effects in the amplification, it is possible to achieve optimal cloning of the qubit encoded in the single-photon polarization. Interestingly, this kind of cloning process does not require any optical nonlinearity to achieve the desired transfer of polarization from the one-photon input to the multi-photon output. Instead, the phase information needed to clone a quantum coherent...
state is obtained explicitly in the form of an optical measurement of field coherence, and the amplification is performed by adding the desired light field amplitude using linear interference between the transmitted light and an appropriately modulated strong laser field. During this process, photon number is not preserved, and the quantum information is transferred from the input photon to the output photon by quantum coherences between states of different photon number. It is thus possible to exploit fundamental aspects of the wave–particle dualism in order to manipulate the discrete polarization statistics of photons through continuous variable operations.

2. Theory of homodyne detection for a single-photon input

Figure 1 shows the schematic setup of the proposed optimal cloning machine. The centre-piece is the beam splitter of reflectivity $R$ that splits the single-photon input into two fields. The quadrature components of the reflected field are then measured by homodyne detection, and a coherent feedback is applied to displace the field amplitude of the transmitted field by $f_R$ times the measurement result $\vec{\beta} = (\beta_H, \beta_V)$. In principle, this setup corresponds to the setups for noiseless amplification [16], compensation of beam splitter losses [15], and Gaussian state cloning [17]. However, in order to handle polarization states, the present setup has to amplify a total of four quadrature components, corresponding to a pair of two orthogonal polarization modes $\hat{a}_H$ and $\hat{a}_V$. For the following discussion, it will be most convenient to define the input state in terms of the creation operators $\hat{a}_H^\dagger$ and $\hat{a}_V^\dagger$ of these two polarization modes, since these

**Figure 1.** Schematic setup of the optimal cloning machine. The one-photon input state $|\psi_{in}\rangle$ is split at a beam splitter of reflectivity $R$. The reflected part is split once more to allow the simultaneous uncertainty limited measurement of the four quadrature components $\hat{x}_H, \hat{x}_V, \hat{y}_H$ and $\hat{y}_V$ by homodyne detection. The measurement result is then transmitted to an optical modulation setup that displaces the transmitted field amplitudes by a feedback of $f_R$ times the measured amplitudes.
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operators also represent the complex conjugate field amplitudes permitting a particularly simple description of linear optics operations on the photon number states. The unknown polarization state of the input photon can then be written in the photon number basis of the two-mode field as

\[ |\psi_{\text{in}}\rangle_a = c_H|0; 1\rangle_H V + c_V|1; 0\rangle_H V = (c_H \hat{a}_H^\dagger + c_V \hat{a}_V^\dagger)|0; 0\rangle_H V. \]  

(1)

The quantum information encoded in this state is expressed by the probability amplitudes \(c_H\) and \(c_V\) of the horizontally (H) and vertically (V) polarized one-photon states.

The effect of the beam splitter on this state can be obtained by transforming the input modes into a coherent superposition of transmitted modes \(\hat{a}_i\) and reflected modes \(\hat{b}_i\),

\[ \hat{U}_{\text{BS}}|\psi_{\text{in}}\rangle_a \otimes |0; 0\rangle_b = (\sqrt{1 - R(c_H \hat{a}_H^\dagger + c_V \hat{a}_V^\dagger)} + \sqrt{R(c_H \hat{b}_H^\dagger + c_V \hat{b}_V^\dagger)})|0; 0\rangle_a \otimes |0; 0\rangle_b. \]  

(2)

The quantum information is now distributed between the reflected modes and the transmitted modes in the form of an entanglement between the modes.

As indicated by figure 1, the reflected mode is then split into equal parts at a beam splitter of reflectivity \(1/2\), and polarization sensitive homodyne detection is applied to both parts to obtain the two quadrature components of the complex amplitudes \(\beta_H = x_H + iy_H\) and \(\beta_V = x_V + iy_V\). The continuous variable measurement of the reflected beam then projects the field in the transmitted beam into a superposition of vacuum and one-photon components, resulting in a corresponding superposition in the transmitted modes. Note that experimentally, this is similar to the recently demonstrated preparation of photonic qubits by field quadrature noise measurements [19]. The main technical difference of our procedure is that we use an additional beam splitter to achieve an uncertainty limited simultaneous measurement of both quadrature components, \(x_H/V\) and \(y_H/V\). In terms of quantum measurement theory, this kind of measurement projects a general input state in the reflected modes \(\hat{b}_i\) onto a coherent field state \(|\beta_H; \beta_V\rangle_{HV}\) where the amplitudes \(\beta_H\) and \(\beta_V\) define the measurement outcome [15]. The properly normalized positive operator valued measure of this measurement reads

\[ |P(\vec{\beta})\rangle = \frac{1}{\pi} |\beta_H; \beta_V\rangle_{HV} \text{ with } \int d^4 \vec{\beta} |P(\vec{\beta})\rangle \langle P(\vec{\beta})| = \hat{1}. \]  

(3)

The conditional output state of the transmitted field modes \(\hat{a}_i\) after the measurement is then given by

\[ \sqrt{p(\vec{\beta})}|\psi(\vec{\beta})\rangle_a = \frac{1}{\pi_b} (\beta_H; \beta_V|\hat{U}_{\text{BS}}|\psi_{\text{in}}\rangle_a. \]  

\[ = \frac{1}{\pi} e^{-|\vec{\beta}|^2/2}(\sqrt{1 - R(c_H \hat{a}_H^\dagger + c_V \hat{a}_V^\dagger)} + \sqrt{R(\beta_H^* c_H + \beta_V^* c_V)})|0; 0\rangle_a. \]  

(4)

where \(p(\vec{\beta})\) is the probability of obtaining the measurement result \(\vec{\beta}\). This state is a coherent superposition of a single-photon state with the original input polarization and a vacuum component, where the quantum coherence between the vacuum and the single-photon states is defined by the relation between the input state amplitudes \(c_H, c_V\) and the measurement results \(\beta_H, \beta_V\).
3. Coherent feedback and optimized gain

It is now possible to modify the output state by coherently adding field amplitudes of $f\beta_H$ and $f\beta_V$ to the polarization components of the output field. As demonstrated in [16, 17], this kind of field addition can be achieved in a straightforward manner by interfering the output field and an appropriately modulated laser beam at a highly reflective beam splitter. However, it may be worth noting that these experiments were performed in continuous wave operation, while a single-photon state must be defined in terms of a finite pulse shape [10]–[14]. In our case, the physical system is thus defined by a single optical pulse, and the conditional displacement must be timed to act on the same pulse from which the measurement data was obtained. This identity of the measured pulse with the output pulse is also the reason why we refer to the conditional displacement as a feedback, in contrast to the terminology used e.g. in [16], where the term ‘feedforward’ is used to indicate the position of the displacement in the continuous beam.

Theoretically, the effects of the measurement on the state in the optical pulse is described by equation (4), and the feedback conditioned by the measurement result can be described by a unitary displacement operator $\hat{D}(f\vec{\beta})$ acting on this conditional output state. (For a detailed discussion of the displacement operator, see e.g. [20].) In order to separate the measurement noise from the input polarization, it is convenient to exchange the ordering of the field operators $\hat{a}_H^\dagger$ and $\hat{a}_V^\dagger$ and the displacement operator $\hat{D}(f\vec{\beta})$ using the following relations,

$$\hat{D}(f\beta_H, f\beta_V)(\hat{a}_H^\dagger f\beta_H^* + \hat{a}_V^\dagger f\beta_V^*) = \hat{a}_H^\dagger \hat{D}(f\beta_H, f\beta_V)(\hat{a}_H^\dagger f\beta_H^* + \hat{a}_V^\dagger f\beta_V^*),$$

$$\hat{D}(f\beta_H, f\beta_V)(\hat{a}_H^\dagger f\beta_H^* + \hat{a}_V^\dagger f\beta_V^*) = \hat{a}_V^\dagger \hat{D}(f\beta_H, f\beta_V).$$

(5)

Comparison with equation (4) then shows that a special feedback condition exists where the displacement can eliminate the products of the input state amplitudes $c_H/V$ and the measurement results $\beta_H/V$. Specifically, the feedback compensated output state for a feedback factor of $f_R = \sqrt{R/(1-R)}$ reads

$$\sqrt{p(\beta)} \hat{D}(f_R\beta_H, f_R\beta_V)|\psi(\beta)\rangle_a = \frac{1}{\pi} e^{-|\vec{\beta}|^2/2} \sqrt{1 - R} (c_H \hat{a}_H^\dagger + c_V \hat{a}_V^\dagger) \hat{D}(f_R\beta_H, f_R\beta_V)|0; 0\rangle_a.$$

(6)

The output state at this special feedback condition is therefore described by the action of the creation operator of the original input photon on a coherent state. As the analogy to cloning by photon bunching [9] suggests, such an application of the single-photon creation operator to an otherwise random state describes an optimal cloning process. For the single-photon input, the unique feedback condition $f_R = \sqrt{R/(1-R)}$ thus converts the beam splitter attenuation into an optimal cloning process.

As first shown in [16], the reason for the existence of the optimal feedback condition can be explained in terms of the quantum noise in the linear field amplification realized by the setup shown in figure 1. Specifically, the feedback condition $f_R = \sqrt{R/(1-R)}$ exactly compensates the effects of the vacuum noise entering at the beam splitter of reflectivity $R$, leaving only the uncertainty limited noise caused by the simultaneous measurement of both quadrature components. Since the setup shown in figure 1 can also be used to amplify and clone coherent states as demonstrated in [17], it may be instructive to express the optimal feedback condition in terms of the field gain that would be obtained for such coherent input fields. If the input was a coherent state with an average amplitude of $\alpha$, the average measurement result $\vec{\beta}$ of the reflected light would be equal to $\sqrt{R}\alpha$ and the feedback would add an average amplitude of $f\sqrt{R}\alpha$ to
the transmitted amplitude of $\sqrt{1 - R}\alpha$, for a total amplitude of $g\alpha$, where the gain factor of the amplification is $g = f\sqrt{R} + \sqrt{1 - R}$. The special feedback condition $fR$ thus corresponds to a gain factor of

$$g_R = fR\sqrt{R} + \sqrt{1 - R} = 1/\sqrt{1 - R}.$$  \hspace{1cm} (7)

Interestingly, this result indicates that optimal cloning is achieved when the gain is exactly the inverse of the attenuation suffered by the transmitted amplitude at the beam splitter.

4. Output density matrix of the optimal cloning process

To show that the effect of this beam splitter amplification on a single photon input is indeed an optimal cloning process, it is necessary to consider the output statistics averaged over all measurement results $\vec{\beta}$. The density operator of the output state is given by

$$\hat{\rho}(\text{out}) = \int d^4\vec{\beta} \ p(\vec{\beta}) \ \hat{D}(fR\vec{\beta}_H, fR\vec{\beta}_V)|\psi(\vec{\beta})\rangle\langle \psi(\vec{\beta})| \hat{D}^\dagger(fR\vec{\beta}_H, fR\vec{\beta}_V) = (1 - R)(c_H^{\dagger}\hat{a}_H + c_V^{\dagger}\hat{a}_V) \ \hat{\eta}_R (c_H^*\hat{a}_H + c_V^*\hat{a}_V),$$ \hspace{1cm} (8)

where the operator $\hat{\eta}_R$ is the density operator of a thermal light field state with an average photon number of $f^2R = R/(1 - R)$ in each mode. The cloning process is thus described by the application of the input photon creation operator to a completely unpolarized light field state. It is now possible to separate $\hat{\rho}(\text{out})$ into contributions with different output photon number $N$,

$$\hat{\rho}(\text{out}) = \sum_{N=1}^\infty P(N) (c_H^{\dagger}\hat{a}_H + c_V^{\dagger}\hat{a}_V) \ \hat{C}_N (c_H^*\hat{a}_H + c_V^*\hat{a}_V),$$ \hspace{1cm} (9)

where $P(N)$ is the probability of an $N$-photon output, and $\hat{C}_N$ is the properly normalized operator of the completely unpolarized $(N - 1)$-photon state before the application of the input photon creation operator,

$$P(N) = \frac{(1 - R)^3}{2R} R^N N(N + 1),$$ \hspace{1cm} (10)

$$\hat{C}_N = \frac{2}{N(N + 1)} \sum_{n=1}^N |n - 1; N - n\rangle \langle n - 1; N - n|. $$ \hspace{1cm} (11)

The average number of clones can be controlled by varying the reflectivity $R$ of the beam splitter, with high reflectivities generating large numbers of clones and low reflectivities generating only a few clones. It is thus far easier to increase the number of clones than in cloning methods relying on parametric downconversion, where it is rather difficult to increase the parametric gain [21]. For practical purposes, however, it may be desirable to keep the cloning probabilities low, since the quantum efficiency of photon detection is usually limited, and the only way to ensure that the detection of $N$ photons really corresponds to $N$ output photons is to keep the probability of generating $N + 1$ photons much lower than the probability for $N$ photons. Equation (10) is therefore essential for the optimal choice of $R$ in an experiment with limited detector efficiencies.
Using the operator $\hat{C}_N$, it is now possible to determine the output statistics of the $1 \rightarrow N$ photon cloning process. The normalized density matrix $\hat{\rho}_N$ of the $N$-photon output reads

$$\hat{\rho}_N = (c_H \hat{a}_H^\dagger + c_V \hat{a}_V^\dagger) \hat{C}_N (c_H^* \hat{a}_H + c_V^* \hat{a}_V). \quad (12)$$

This output is a mixture of photon number states with $n$ photons in the correct input polarization and $N - n$ photons in the opposite polarization. The statistical weight of each state is determined by the factor of $n$ introduced by the application of the creation operator of the input photon to both sides of the unpolarized operator $\hat{C}_N$. The normalized probability distribution $P(n|N)$ of the number of correctly polarized photons $n$ among $N$ output photons thus reads

$$P(n|N) = \frac{2n}{N(N+1)}. \quad (13)$$

The fidelity of the $1 \rightarrow N$ photon cloning process is then given by the ratio between the average photon number in the input polarization and the total output photon number,

$$F_{1 \rightarrow N} = \sum_{n=0}^{N} P(n|N) \frac{n}{N} = \frac{2N + 1}{3N}. \quad (14)$$

This is the optimal fidelity for $1 \rightarrow N$ cloning [22]. Thus, the coherent feedback setup shown in figure 1 is indeed an optimal cloning machine.

The analysis above assumed precise homodyne detection and field displacements. In realistic implementations, it will be necessary to take into account additional errors. In particular, nonunit quantum efficiency and mode matching may introduce errors in the linear amplification scheme, as discussed in some detail in [18]. While a detailed analysis is beyond the scope of the present paper, it may be worthwhile to consider the possible effects of such errors on the $N$-photon polarization observed in the output. Assuming that the errors are not polarization sensitive, we may assume that they can be represented by a ‘white noise’ background of equal probability, as given by the density operator $\hat{W}_N = \hat{C}_{N+1}$. The actual cloning fidelity achieved will then be a weighted average of the optimal cloning fidelity and the ‘white noise’ fidelity of $1/2$.

5. Information and noise in the coherent cloning process

As the discussion above has shown, it is in principle possible to realize optimal cloning of a single-photon polarization state by applying quantum measurements to the continuous field variables. The field measurement projects the transmitted field into a coherent superposition of vacuum and single-photon components as given by equation (4). By choosing an optimal feedback gain of $g_R = 1/\sqrt{1 - \tilde{R}}$, this superposition of vacuum and single-photon component can be converted into a superposition of optimally cloned $N$-photon outputs, as given by equation (6). Remarkably, this manipulation of photon polarization states is achieved entirely by continuous variable operations, making use of the correspondence between the quantum coherence of the single photon state and the (classical) optical coherence of the field.

It may also be worth noting that, due to the formal equivalence of continuous variable teleportation errors and beam splitter losses [15], a closely related optimal cloning process can be implemented by the continuous variable teleportation of single-photon states [23–25].
In this case, the optimal gain condition depends on the squeezed state entanglement, with low entanglement requiring a correspondingly higher gain to achieve optimal cloning. At a teleportation gain of $g = 1$, the cloning process is not optimal and the teleportation errors in the $N$ photon outputs will be greater than the minimal cloning errors [26]. In both cases, the essential feature of the cloning process is that continuous variable measurements and field displacements are used to clone the polarization states of individual input photons.

Besides demonstrating the potential usefulness of continuous variable operations for the processing of photon polarization qubits, the other significant feature that distinguishes the present cloning scheme from previous proposals for the cloning of photon polarization is the use of projective measurements to implement optimal cloning. It may be interesting to note that the cloning process can be optimized despite the partial loss of the original photon at the beam splitter. Obviously, the loss of quantum information due to the partial absorption of the input photon is balanced by the usefulness of the classical measurement information $\vec{\beta}$ for the generation of quantum clones. This balance could also be used to explain the existence of an optimal gain: at $g > 1/\sqrt{1 - R}$, the measurement information obtained is not sufficient to optimize the fidelity of the high number of clones generated, and at $g < 1/\sqrt{1 - R}$, the measurement back action caused by the unnecessary precision of the measurement introduces additional cloning errors. The present cloning methods may thus provide some insights into the relation between classical information and quantum information in quantum cloning processes.

6. Conclusions

In conclusion, we have shown that an amplification of the light field by coherent feedback of the reflection losses at a beam splitter of reflectivity $R$ optimally clones the polarization state of a single-photon input if the feedback induced gain is equal to $g_R = 1/\sqrt{1 - R}$. This cloning method does not require any nonlinear optical elements and multiple clones are easy to obtain. It may therefore be particularly useful for closing the gap between output photon numbers of $N = 2$ and $N \to \infty$. By employing field measurements to manipulate the polarization states of photons, this cloning methods also illustrates the fundamental relation between the continuous field variables and the discrete photon number distributions of the quantized light field. Photon cloning by coherent feedback amplification thus shows how fundamental aspects of the wave–particle dualism of light can be applied to realize quantum information processes.

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