Reversible Decoherence of a Mesoscopic Superposition of Field States

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We describe the blueprint of an experiment in which the decoherence of a mesoscopic superposition of radiation states ("Schrödinger cat") becomes a reversible process. When the high Q cavity containing the Schrödinger cat is coupled to another resonator, the mesoscopic quantum coherence first decays rapidly, then exhibits sharp revivals with the period of the energy exchange between the two cavities. The interpretation of this experiment emphasizes the link between decoherence and complementarity and leads to an illuminating quantitative interpretation of the usual irreversible decoherence phenomenon. [S0031-9007(97)04031-3]

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The absence of macroscopic quantum superpositions is a central issue in our understanding of quantum measurement theory [1]. Unitary quantum evolution predicts that the meter in a measurement process should generally evolve into a quantum superposition of outcomes. Schrödinger [2] has vividly illustrated this problem by replacing the meter by a cat, whose life is dependent upon the fate of a radioactive atom. This situation leads to paradoxical superpositions of "dead" and "alive" cat "states". By reference to this metaphor, macroscopic or mesoscopic quantum superposition states are often called "Schrödinger cats".

Recent theoretical advances have stressed the role of the environment in the solution of this paradox [3,4]. Macroscopic systems are strongly coupled to large reservoirs with many degrees of freedom. In any realistic model, this coupling very rapidly blurs quantum superpositions, changing them into mere statistical mixtures. This "decoherence" becomes increasingly faster with the size of the system. It can be easily interpreted in terms of complementarity. Due to the coupling with the reservoir, information about the system continuously leaks into the environment, until the states of the reservoir correlated to the different states of the system become orthogonal. The system's reduced density matrix turns continuously into a statistical mixture. In this respect, the decoherence process is viewed as a succession of uncontrolled and unread "measurements" of the system by its environment [5]. The time scale of this information gathering process decreases when the size of the system increases. A recent experiment [6] using circular Rydberg atoms [7] and a superconducting cavity has explored for the first time the dynamics of this process on a Schrödinger cat made of a superposition of two coherent fields [8] with different phases (a "phase cat" [9]).

Usually, decoherence blends the concepts of complementarity (leakage of information) and irreversibility (a consequence of the large size of the environment). These two aspects are not necessarily linked, however, and it is possible, by controlling the "size" of the reservoir, to design situations in which the coherence of a Schrödinger cat, instead of decaying irreversibly, "collapses" and "revives" periodically. We propose in this Letter an experimental scheme derived from [6] to test decoherence in such a well-controlled environment. By coupling the cavity containing a phase cat to another resonator playing the role of a "single mode reservoir," a reversible exchange of energy between the two cavities can be achieved and a "reversible decoherence" process could be observed. The evolution of the cat's coherence is revealed by a quantum interference signal involving two "paths" corresponding to the interaction of an atom with the two field phase components. This interference signal disappears, at the beginning of the system's evolution, as soon as the single mode "reservoir" contains a field carrying unambiguous information about the cavity field phase. The interference is restored after one period of the energy exchange between the cavities, when the reservoir is empty again. Successive collapses and revivals of the "cat's" coherence could in principle be observed.

The situation is reminiscent of "quantum eraser" experiments in which a particle undergoes an interference process [10]. Interferences disappear when information about the particle's path is available, encoded in a complementary microsystem. They can be restored by manipulating this system in a way which "erases" the "which path" information. In the case considered in this Letter, however, the loss and revivals of the coherence are dynamical processes, occuring on time scales becoming shorter when the size of the field in the cavity is increased, an essential feature of mesoscopic systems. The proposed experiment would demonstrate the essential role of complementarity in the decoherence process. When the cavity is coupled to more and more "reservoir" resonators, we recover the results of the standard irreversible decoherence model (cavity interacting with a bath of harmonic oscillators). This new approach is simpler and more illuminating than other standard computations [11,12].

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The proposed scheme is sketched on Fig. 1. A single circular Rydberg atom is used to prepare, at time $t = 0$, a phase cat in the cavity C_0, which is coupled to another resonant, initially empty, cavity C_1. The frequency of the energy exchange between C_0 and C_1 is $\Omega_c/2\pi$. It could be tuned at will by adjusting the cavities coupling through a superconducting waveguide. We will neglect relaxation processes for the cavities, as well as for the atoms, during the experiment duration. This is realistic for experimentally achievable cavity quality factors of a few 10^9, corresponding to photon lifetimes T_R of the order of a few ms. The principle of the phase cat preparation is described in detail in [6]. Let us note that the coupling between C_0 and C_1 plays no role during the cat generation, provided the preparation time is much shorter than Ω_c^{-1}. This time is of the order of the atomic transit time through the cavity, t_i, about 20 μs. Since it is much smaller than T_R, Ω_c can be chosen such that the inequalities $t_i \ll \Omega_c^{-1} \ll T_R$ are simultaneously fulfilled, which we assume in the following.

The "cat" generation makes use of the dispersive, nonresonant coupling of a single circular Rydberg atom to the cavity mode, in which a small coherent field $|\alpha_0\rangle$ is initially prepared by a pulsed classical microwave source S [6]. The average photon number $n = |\alpha_0|^2$ is typically varied between 0 and 10. The atom, effusing from an oven O, velocity selected in zone V, is excited into a circular Rydberg state in zone B. It is prepared, before entering C_0, in a superposition of two circular Rydberg states e and g (principal quantum numbers 51 and 50, respectively) by a classical microwave field applied in zone R_1. This field performs the transformation $|e\rangle \rightarrow (|e\rangle + |g\rangle)/\sqrt{2}$. The detuning δ between the cavity frequency $\omega/2\pi$ and the $e \rightarrow g$ atomic transition frequency at 51.099 GHz is large enough to avoid any energy transfer between the atom and C_0. The atom behaves thus as a piece of transparent dielectric, with an index of refraction large enough to modify appreciably the phase of the cavity field during the atom transit time t_i. When in level e, the atom changes the cavity field phase by an angle ϕ, yielding a cavity state $|\alpha_0 \exp i\phi\rangle$. An atom in level g leaves in C_0 the state $|\alpha_0 \exp -i\phi\rangle$. An atomic superposition of e and g prepares a field state involving a superposition of these two phase components. The phase ϕ reaches high values in the experiment (0.7 rad for $\delta/2\pi = 100$ kHz). After interaction with the cavity, the atomic states are mixed again in a second classical microwave zone (R_2) performing the transformations $|e\rangle \rightarrow (|e\rangle + |g\rangle)/\sqrt{2}$ and $|g\rangle \rightarrow (-|e\rangle + |g\rangle)/\sqrt{2}$. We choose here for simplicity the same phase for the classical fields in R_1 and R_2, both fed by source S'. The atom is finally detected by field ionization counters D_e and D_g, either in state e or in state g. Since R_2 erases any information on the atomic state in C_0, the detection projects the cavity state onto the phase cat state:

$$|\Psi_c\rangle = \frac{1}{\sqrt{2}}\left(|\alpha_0 e^{i\phi}\rangle \pm |\alpha_0 e^{-i\phi}\rangle\right), \quad (1)$$

where the $+$ sign applies for a detection in g, and the $-$ sign for a detection in e [13]. We have neglected in the normalization factor the overlap between the two coherent components, assuming that $|\alpha_0| \gg 1$ and $\phi \neq 0$ or $\pi/2$.

The phase cat left in the cavity is probed after a fixed delay τ by a second atom crossing the setup. Let us assume first that C_1 is not coupled to C_0 ($\Omega_c = 0$). Being also in a quantum superposition of two states, the second atom splits again the two components of the phase cat, adding its own dephasing to the first atom's one [6]. The final field state therefore contains four phase components. Two of them overlap at zero phase. They correspond to the two quantum paths where the second atom undoes the phase shift of the first one (atoms cross C_0 in the configurations e, g or g, e). Since the atomic levels are mixed in R_2, these two paths are indistinguishable, resulting in a quantum interference. It yields correlations between the two detected atomic states. The correlation signal η is the difference between the conditional probability Π_{ee} to detect the second atom in e provided the first was in e, and the conditional probability Π_{ge} to detect the second in e if the first was in g. In a calculation generalizing [11], the correlation is found to be equal to half the real part of the overlap of the two field components at zero phase: $\eta = 1/2\,\mathrm{Re}\langle\alpha_0 \,|\, \alpha_0\rangle = 1/2$. This analysis neglects the components with phases $\pm 2\phi$ which correspond to the atomic paths e, e or g, g in C_0. These two field components do not overlap, provided again that $|\alpha_0| \gg 1$ and $\phi \neq 0, \pi/2$, a condition assumed to be true in the following [relaxing this condition, and hence also the assumption that the two components in Eq. (1) do not overlap, adds only minor algebraic complications].

Let us now take into account the coupling between C_0 and C_1 during the delay τ. Consider first the case where C_0 initially contains a single coherent component $|\alpha_0\rangle$. Due to the linearity of the coupling between the two cavities, each of them contains at any time a coherent field: $|\alpha(t)\rangle$ in C_0 and $|\beta(t)\rangle$ in C_1. Since C_1 is initially empty, one gets $\alpha(t) = \alpha_0 \cos(\Omega_c t/2)$ and

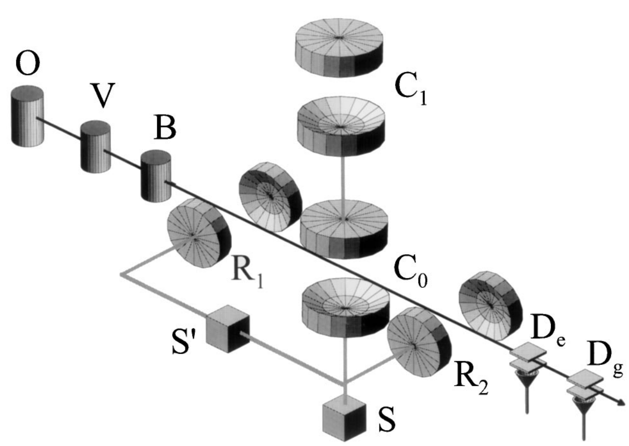

FIG. 1. Sketch of the proposed setup.

$\beta(t) = \alpha_0 \sin(\Omega_c t/2)$ (we make here a proper choice for the field phase origin in the two cavities). The energy is oscillating between C_0 and C_1 at frequency $\Omega_c/2\pi$, a purely classical effect. If a "cat" state $|\Psi_c\rangle$ is initially prepared in C_0, the field in $C_0 + C_1$ becomes the entangled state: $(|\alpha(t)e^{i\phi}\rangle|\beta(t)e^{i\phi}\rangle \pm |\alpha(t)e^{-i\phi}\rangle|\beta(t)e^{-i\phi}\rangle)/\sqrt{2}$.

When the second atom probes the field in C_0, C_1 is not affected (provided again that the "measurement time" is much less than Ω_c^{-1}). The C_0 field again involves a superposition of four phase components. Two of them merge at zero phase. The correlation signal η is proportional to the overlap of the two–cavity–field states corresponding to the two quantum paths giving rise to the interference process. The e, g path corresponds to the final two–cavity field state $|\alpha(\tau)\rangle|\beta(\tau)e^{i\phi}\rangle$ whereas the g, e path corresponds to $|\alpha(\tau)\rangle|\beta(\tau)e^{-i\phi}\rangle$. The η value is then given by

$$\begin{aligned}\eta &= \frac{1}{2}\,\mathrm{Re}\left(\left\langle \alpha_0 \cos\left(\frac{\Omega_c}{2}\tau\right) \middle| \alpha_0 \cos\left(\frac{\Omega_c}{2}\tau\right)\right\rangle \left\langle \alpha_0 \sin\left(\frac{\Omega_c}{2}\tau\right)e^{-i\phi} \middle| \alpha_0 \sin\left(\frac{\Omega_c}{2}\tau\right)e^{i\phi}\right\rangle\right) \\ &= \frac{1}{2}\,e^{-2n\sin^2(\Omega_c\tau/2)\sin^2\phi}\cos[n\sin^2(\Omega_c\tau/2)\sin 2\phi] \end{aligned} \tag{2}$$

(Note that the scalar product of two coherent states $|a\rangle$ and $|b\rangle$ is $\langle a\,|\,b\rangle = e^{-(|a|^2+|b|^2)/2}e^{ba^*}$ [8]). The assumption that the field components left in C_0 do not overlap fails for times τ around π/Ω_c, when all the energy is localized in C_1. At such times, the above expression is not valid. However, since C_0 is empty, η is expected to be zero, which is correctly predicted by the above equation.

The corresponding signal $\eta(\tau)$ is plotted on Fig. 2, for $n = 5$ in the initial field. The delay τ is expressed in units of the energy exchange period, $2\pi/\Omega_c$. The initial coherence is rapidly washed out when a field leaks into C_1. For a time τ much smaller than Ω_c^{-1}, η is $\simeq (1/2)\exp[-n\Omega_c^2(\sin^2\phi)\tau^2/2]$. The time scale of the η decay is therefore $\sqrt{2/n}\,(\Omega_c \sin\phi)^{-1}$. Note that, at variance with the standard decoherence process, this time scale is inversely proportional to the field amplitude and not to the field energy.

The disappearance of the quantum correlation can be understood as a complementarity effect. The field building up in C_1, which plays the role of an environment for C_0, carries away information about the phase of the field in C_0. As soon as there is in C_1 enough information to determine, at least in principle, the phase in C_0, the quantum coherence of the mesoscopic field in C_0 is lost. In more technical words, as soon as the two phase components in C_1 are orthogonal, the reduced density matrix of the C_0 field describes a statistical mixture of the two-phase components. If ϕ is large enough, the phase uncertainty on the "leak field components" in C_1 will be smaller than 2ϕ as soon as the leak field energy corresponds to one or two photons. The time to reach such a leak amplitude is obviously of the order of $1/(\sqrt{n}\,\Omega_c)$. At variance with standard decoherence, the leak is a reversible process: after one or several full periods of the energy exchange, the field returns completely in C_0 with the initial phases components. A revival of the coherence signal is observed around these times (see Fig. 2). With an environment made of a single quantum oscillator, decoherence becomes reversible, in close analogy with the spontaneous emission of an atom in a single undamped mode [14].

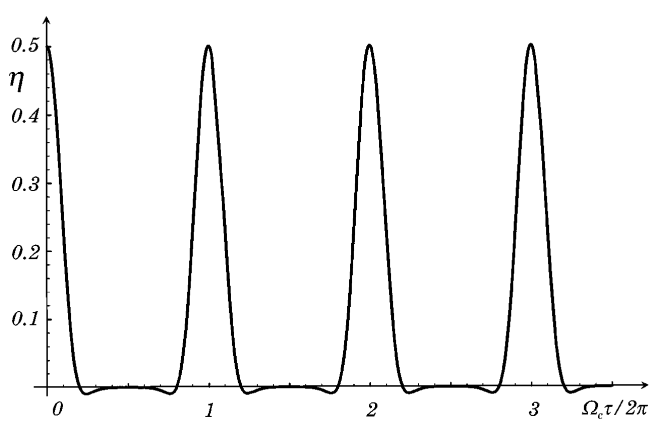

FIG. 2. Two-atom correlation signal $\eta(\tau)$ for an initial field in C_0 containing on the average 5 photons. The delay τ between the atoms is expressed in units of the energy exchange period between C_0 and C_1, $2\pi/\Omega_c$.

This very simple approach to decoherence in terms of complementarity can be used also to recover the usual decoherence theory result for an exponential damping of the energy in C_0 [11,12]. Let us assume that C_0 is now coupled to a large number of cavities $C_1, C_2, \ldots, C_i, \ldots$, with arbitrary frequencies (close to the one of C_0) and arbitrary coupling constants. This models the standard situation of an harmonic oscillator relaxing in a large bath of zero temperature oscillators. An initially coherent field remains coherent, with an energy decreasing exponentially as $\exp(-\gamma t)$ ($\gamma = 1/T_R$). When C_0 initially contains a phase cat, each of its components is correlated in C_i to a small field $|\beta_i(t)e^{i\phi}\rangle$ or $|\beta_i(t)e^{-i\phi}\rangle$. Each cavity carries a very small amount of information on the phase in C_0, but there are quite many of them. By a straightforward generalization of the preceding arguments, the correlation signal obtained for a second atom crossing C_0 at time τ will be

$$\begin{aligned}\eta(\tau) &= \frac{1}{2}\,\mathrm{Re}\prod_i \langle \beta_i(\tau)e^{-i\phi}\,|\,\beta_i(\tau)e^{i\phi}\rangle \\ &= \frac{1}{2}\,\mathrm{Re}\exp\left(-\sum_i |\beta_i(\tau)|^2(1 - e^{-2i\phi})\right). \end{aligned} \tag{3}$$

Due to energy conservation, $\sum_i |\beta_i(\tau)|^2$ is the average number of photons having escaped from C_0 at time τ, i.e.,

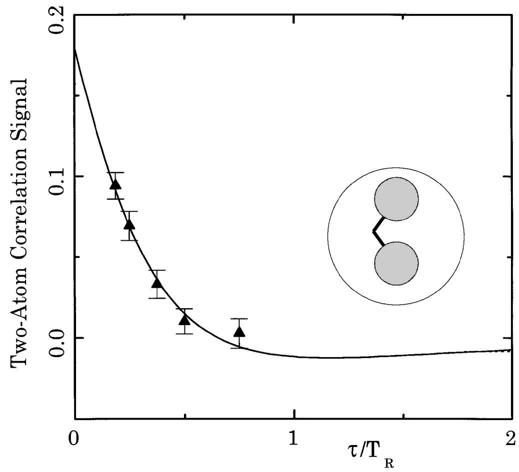

FIG. 3. Two-atom correlation signal $\eta(\tau)$ for an initial 3.3 photons coherent field in C_0, coupled to a bath of harmonic oscillators. The two field phase components left by the first atom are represented in the inset ($\phi = 1$ rad). The delay τ between the atoms is expressed in units of T_R. The points are experimental (see [6]), and the curve results from the theoretical model [Eq. (4)].

$n(1 - e^{-\gamma\tau})$. The correlation signal writes finally

$$\eta(\tau) = \frac{1}{2} e^{-2n(1-e^{-\gamma\tau})\sin^2\phi} \cos[n(1 - e^{-\gamma\tau}) \sin 2\phi] . \tag{4}$$

Note that this signal coincides, in the frame of our approximations, with the result of an exact calculation, adapted from [11]. This decoherence signal is presented in Fig. 3 for an $n = 3.3$ photons field and $\phi = 1$ rad, superimposed to the experimental signal obtained in [6]. The theoretical signal of Eq. (4) has been normalized to a maximum value of 0.18 to account for experimental imperfections (see [6]).

By stressing its connection to complementarity, we have shown that decoherence is not necessarily an irreversible process and we have described a possible experiment to demonstrate the revival of mesoscopic coherences at long times. The main difficulty of this experiment is the realization of a reversible coupling between two identical Fabry-Perot cavities. An improvement of the techniques already used to couple external microwave sources to the cavity in our setup [6] could make it possible. Many variants of this double cavity experiment can also be considered. By sending an atom beam through C_1 as well as C_0, one could, for example, manipulate directly the field in the reservoir cavity and implement various "quantum eraser" schemes [10].

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