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Control of non-Markovian decay and decoherence by measurements and interference

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Abstract: Novel methods are discussed for the state control of atoms coupled to multi-mode reservoirs with non-Markovian spectra:
1) *Excitation decay control*: we point out that the quantum Zeno effect, i.e., inhibition of spontaneous decay by frequent measurements, is observable in open cavities and waveguides using a sequence of evolution-interrupting pulses or randomly-modulated CW fields.
2) *Location-dependent interference of decay channels - nonadiabatic (resonant) control*: We show that the control of populations and coherences of two metastable states is feasible via *resonant* single-photon absorption to an intermediate state, by controlled spontaneous emission in a cavity.
3) *Decoherence control by conditionally interfering parallel evolutions*: We demonstrate that an *arbitrary* internal atomic state can be completely protected from decoherence by interference of its interactions with the reservoir over many different time intervals *in parallel*. Such interference is conditional upon the detection of appropriate atomic-momentum observables. Realization in cavities is suggested.

The rich arsenal of control methods described above can improve the performance of single-atom devices. It can also advance the state-of-the-art of quantum information encoding and processing.

OCIS codes: (270.2500) Fluctuations, relaxations, and noise; (020.0020) Atomic and molecular physics

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1. Introduction

It is our purpose here to demonstrate that effective control of atomic states is possible in the presence of spontaneous relaxation into non-Markovian reservoirs, characterized by non-smooth spectra. Cavities, waveguides and periodic dielectric structures are prime examples of environments that give rise to such reservoirs for electromagnetic (EM) field modes, whereas condensed media or multi-ion traps are their analogs for phonon modes. In cavities, one may employ the following arsenal of means for the control of radiative decay: (i) *mode density control*, which allows the suppression or enhancement of the decay rate in selected spectral ranges; (ii) *location control*, which allows the adjustment of the strength and phase of coupling to a standing-wave field mode of the cavity, by varying the position (relative to field nodes) of atoms or molecules deposited on a thin film; and (iii) *measurement control*, namely, non-selective or conditional measurements (CMs) of an *auxiliary variable* that is correlated to the internal state of the atom.

In what follows, we shall demonstrate how this arsenal can be implemented in several methods we have developed for quantum state control in cavities: (a) population decay control by frequent or continuous measurements, realizing the quantum Zeno effect[1]; (b) complete control of state preparation by location-dependent interference of decay channels[2]; and (c) control of decoherence by CM-induced interference of evolution histories[3]. The aim will be to clarify the essentials of these methods, by means of computer movies and other graphical illustrations, and finally compare their advantages and disadvantages.

2. Quantum Zeno effect: control of excitation decay

The "watchdog" or quantum Zeno effect (QZE) is a spectacular manifestation of the influence of measurements on the evolution of a quantum system. The original QZE prediction has been that irreversible decay of an excited state into a reservoir can be inhibited[4], by repeated interruption of the system-reservoir coupling, which is associated with measurements[5,6]. However, this prediction has not been experimentally verified as yet! Instead, the interruption of Rabi oscillations and analogous forms of *nearly-reversible* evolution has been at the focus of interest[7–14].

We have recently demonstrated[15] that the inhibition of *nearly-exponential* excited-state decay by the QZE in two-level atoms, in the spirit of the original suggestion[4], is amenable to experimental verification in resonators. Although this task has been widely believed to be very difficult, we have shown, by means of our unified theory of spontaneous emission into reservoirs with arbitrary mode-density spectra[16], that several realizable configurations based on two-level emitters in cavities or in waveguides are in fact adequate for QZE observation. We have now developed a more comprehensive view of the possibilities of excited-state decay by QZE[1]. Here we wish to demonstrate, by means of computer movies and illustrations, that QZE is indeed achievable by repeated or continuous measurements of the excited state, in various non-Markovian reservoirs.

2.1 QZE by frequent impulsive measurements

Consider an initially excited two-level atom coupled to an *arbitrary* density-of-modes (DOM) spectrum $\rho(\omega)$ of the electromagnetic field in the vacuum state. At time τ its evolution is interrupted by a short optical pulse, which serves as an impulsive quantum measurement[7–14]. Its role is *to break the evolution coherence*, by transferring the populations of the excited state $|e\rangle$ to an auxiliary state which then decays back to $|e\rangle$ *incoherently.*

The atomic response, i.e., the emission rate into this reservoir at frequency ω, which is $|g(\omega)|^2\rho(\omega)$, $\hbar g(\omega)$ being the field-atom coupling energy, can be divided into two parts,

$$G(\omega) = G_s(\omega) + G_b(\omega). \tag{1}$$

Here $G_s(\omega)$ stands for the sharply-varying (nearly-singular) part of the DOM distribution, associated with narrow cavity-mode lines or with the frequency cutoff in waveguides or photonic band edges. The complementary part $G_b(\omega)$ stands for the broad portion of the DOM distribution (the "background" modes), which always coincides with the free-space DOM $\rho(\omega) \sim \omega^2$ at frequencies well above the sharp spectral features. In an open structure (see below), $G_b(\omega)$ represents the atom coupling to the unconfined free-space modes.

We cast the excited-state amplitude in the form $\alpha_e(\tau)e^{-i\omega_a\tau}$, where ω_a is the atomic resonance frequency. Restricting ourselves to sufficiently short interruption intervals τ such that $\alpha_e(\tau) \simeq 1$, yet long enough to allow the rotating wave approximation, we obtain

$$\alpha_e(\tau) \simeq 1 - \int_0^\tau dt(\tau - t)\Phi_s(t)e^{i\Delta t} - \gamma_b\tau/2, \tag{2}$$

where

$$\Phi_s(t) = \int_0^\infty d\omega G_s(\omega)e^{-i(\omega-\omega_s)t}. \tag{3}$$

$\Delta = \omega_a - \omega_s$ is the detuning of the atomic resonance from the peak ω_s of $G_s(\omega)$, and $\gamma_b = 2\pi G_b(\omega_a)$ is the effective rate of spontaneous emission into the background modes.

To first order in the atom-field interaction, the excited state probability after n interruptions (measurements), $W(t = n\tau) = |\alpha_e(\tau)|^{2n}$, can be written as

$$W(t = n\tau) \approx [2\mathrm{Re}\alpha_e(\tau) - 1]^n \approx e^{-\kappa t}, \tag{4}$$

where

$$\kappa = 2\mathrm{Re}[1 - \alpha_e(\tau)]/\tau. \tag{5}$$

In most structures γ_b is comparable to the free-space decay rate γ_f and gives rise to an *exponential* decay factor in the excited state probability, regardless of how short τ is, i.e.,

$$\kappa = \kappa_s + \gamma_b, \tag{6}$$

where κ_s is the contribution to κ from the sharply-varying modes.

Thus the background-DOM effect cannot be modified by QZE. Only the sharply-varying DOM contribution κ_s allows for QZE, provided that

$$\kappa_s = (2/\tau)\mathrm{Re}\int_0^\tau dt(\tau - t)\Phi_s(t)e^{i\Delta t} \tag{7}$$

decreases with τ for sufficiently short τ. This essentially means that the correlation (or memory) time of the field reservoir is longer (or, equivalently, $\Phi_s(t)$ falls off slower) than the chosen interruption interval τ.

2.1 .1 Application to a Lorentzian line

The simplest application of the above analysis is to the case of a two-level atom coupled to a near-resonant Lorentzian line centered at ω_s, characterizing a high-Q cavity mode[15]. In this case, $G_s(\omega) = g_s^2\Gamma_s/\{\pi[\Gamma_s^2+(\omega-\omega_s)^2]\}$, where g_s is the resonant coupling strength and Γ_s is the linewidth (Fig. 1). In the short-time approximation, taking into account that the Fourier transform of the Lorentzian $G_s(\omega)$ is $\Phi_s(t) = g_s^2e^{-\Gamma_s t}$, Eq. (2) yields (without the background-modes contribution)

$$\alpha_e(\tau) \approx 1 - \frac{g_s^2}{\Gamma_s - i\Delta}\left[\tau + \frac{e^{(i\Delta-\Gamma_s)\tau} - 1}{\Gamma_s - i\Delta}\right]. \tag{8}$$

The QZE condition is then

$$\tau \ll (\Gamma_s + |\Delta|)^{-1}, g_s^{-1}. \tag{9}$$

On resonance, when $\Delta = 0$, Eq. (8) yields

$$\kappa = \kappa_s + \gamma_b, \quad \kappa_s = g_s^2\tau. \tag{10}$$

Only the κ_s term decreases with τ, indicating the QZE inhibition of the nearly-exponential decay into the Lorentzian field reservoir as $\tau \to 0$. Since Γ_s has dropped out of Eq. (10),

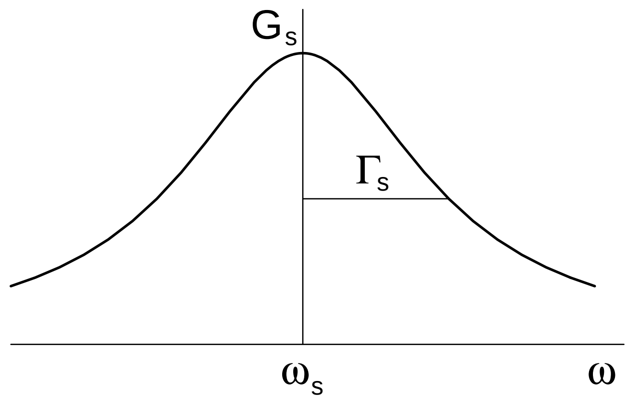

Figure 1: Cavity mode with Lorentzian lineshape

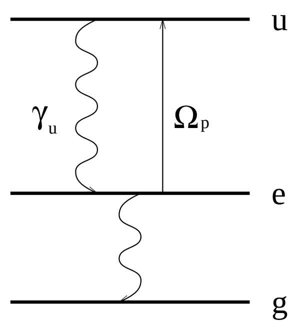

Figure 2: The level scheme

the decay rate κ is the *same* for both strong-coupling ($g_s > \Gamma_s$) and weak-coupling ($g_s \ll \Gamma_s$) regimes. Physically, this comes about since for $\tau \ll g_s^{-1}$ the energy uncertainty of the emitted photon is too large to distinguish between reversible and irreversible evolutions.

The evolution inhibition, however, has rather different meaning for the two regimes. In the weak-coupling regime, where, in the absence of the external control, the excited-state population decays nearly exponentially at the rate $g_s^2/\Gamma_s + \gamma_b$ (at $\Delta = 0$), one can speak about the inhibition of irreversible decay, in the spirit of the original QZE prediction[4]. By contrast, in the strong-coupling regime in the absence of interruptions (measurements), the excited-state population undergoes damped Rabi oscillations at the frequency $2g_s$. In this case, the QZE slows down the evolution during the first Rabi half-cycle ($0 \leq t \leq \pi/2g_s^{-1}$), the evolution on the whole becoming irreversible.

A possible realization of this scheme is as follows. Within an open cavity the atoms repeatedly interact with a pump laser, which is resonant with the $|e\rangle \rightarrow |u\rangle$ transition frequency. The resulting $|e\rangle \rightarrow |g\rangle$ fluorescence rate is collected and monitored as a function of the pulse repetition rate $1/\tau$. Each short, intense pump pulse of duration t_p and Rabi frequency Ω_p is followed by spontaneous decay from $|u\rangle$ back to $|e\rangle$, at a rate γ_u, so as to *destroy the coherence* of the system evolution, on the one hand, and *reshuffle the entire population* from $|e\rangle$ to $|u\rangle$ and back, on the other hand (Fig. 2). The demand that the interval between measurements significantly exceed the measurement time,

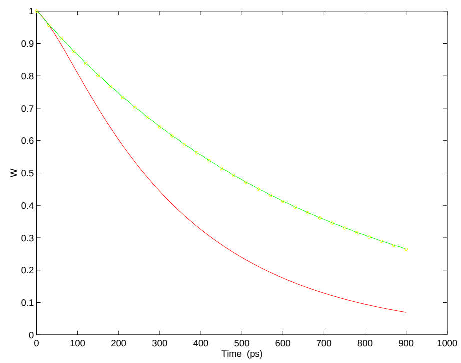

Figure 3: Movie A: Evolution of excited-state population W in two-level atom coupled to cavity mode with Lorentzian lineshape on resonance, $(\Delta = 0)$: red curve – uninterrupted decay in cavity with $F \equiv (1-R)^{-2} = 10^4$, L=15 cm, and f=0.02; green curve – interrupted evolution at intervals $\tau = 3 \times 10^{-8}$ s, yellow dots denote the interruption moments. Here $\gamma_b \simeq \gamma_f = 10^6$ s^{-1};

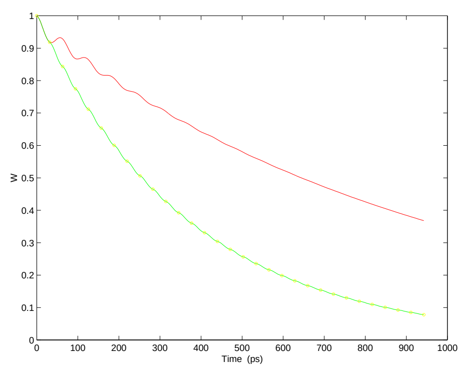

Figure 4: Movie B: Idem, for detuning $\Delta = 10^8 \mathrm{s}^{-1}$ and $F = 10^5$.

yields the inequality $\tau \gg t_p$. The above inequality can be reduced to the requirement $\tau \gg \gamma_u^{-1}$ if the "measurements" are performed with π pulses: $\Omega_p t_p = \pi,\ t_p \ll \gamma_u^{-1}$. This calls for choosing a $|u\rangle \to |e\rangle$ transition with a much shorter radiative lifetime than that of $|e\rangle \to |g\rangle$.

Movie A, describing the QZE for a Lorentz line on resonance $(\Delta = 0)$, has been programmed for feasible cavity parameters: $\Gamma_s = (1-R)c/L,\ g_s = \sqrt{cf\gamma_f/(2L)},\ \gamma_b = (1-f)\gamma_f$, where R is the geometric-mean reflectivity of the two mirrors, f is the fractional solid angle (normalized to 4π) subtended by the confocal cavity, and L is the cavity length. It shows, that the population of $|e\rangle$ decays nearly-exponentially well within interruption intervals τ, but when those intervals become too short, there is significant inhibition of the decay. Movie B shows the effect of the detuning $\Delta = \omega_a - \omega_s$ on the decay: The decay now becomes oscillatory. The interruptions now *enhance* the decay, the degree of enhancement depends on the phase between interruptions.

2.1 .2 Application to mode distributions with cutoff

We can extend the above analysis to DOM distributions characterized by *a cutoff frequency*, as in a waveguide, a photonic band edge or *a phonon reservoir* (with Debye cutoff). A specific model for the spectral response of a DOM distribution with a cutoff

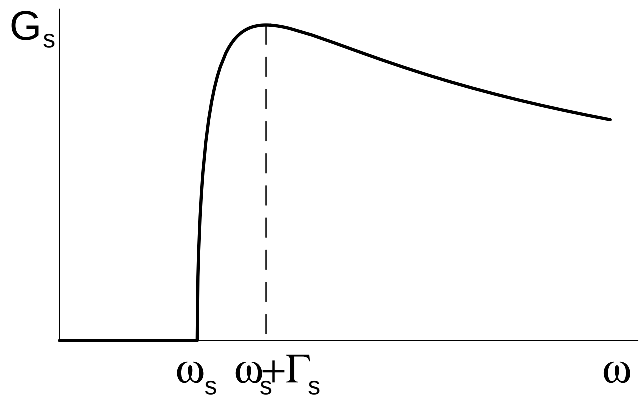

Figure 5: DOM with cutoff

is represented by[16]

$$G_s(\omega) = [C\sqrt{\omega - \omega_s}/(\omega - \omega_s + \Gamma_s)]\Theta(\omega - \omega_s), \tag{11}$$

where ω_s is the cutoff (or band-edge) frequency, Γ_s is the cutoff-smoothing parameter, C is the strength of the coupling of the atomic dipole to this reservoir, and $\Theta()$ is the Heaviside step function (Fig. 5). Upon computing the Fourier transform of Eq. (11), we find from Eqs. (3),(7) that the QZE condition is

$$\tau \ll \min\{\Gamma_s^{-1}, |\Delta|^{-1}, C^{-2/3}\}. \tag{12}$$

Under this condition, Eq. (7) yields

$$\kappa_s = (2^{5/2}\pi^{1/2}/3)C\tau^{1/2}. \tag{13}$$

As seen in our movie C, the QZE is now *less pronounced* as compared to the Lorentzian-line case (movie A). This case is realizable for an active dipole layer embedded in a dielectric waveguide, using a level scheme similar to that of Fig. 2. The same evolution obtains for a state relaxing vibrationally into a phonon reservoir, if the vibrational transition is near the Debye cutoff of the reservoir.

2.2 QZE by continuous measurement: noisy-field dephasing

Instead of disrupting the coherence of the evolution by a sequence of "impulsive" measurements, i.e., short π-pulses, we can achieve this goal by *noisy-field dephasing* of $\alpha_e(t)$: Random ac-Stark shifts by an off-resonant intensity-fluctuating field result in the replacement of Eq. (7) by (Fig. 7)

$$\kappa_s = \int G_s(\Delta + \omega_a)F(\Delta)d\Delta, \tag{14}$$

Here the sharply varying spectral response $G_s(\Delta + \omega_a)$ [Eq. (1)] replaces the Fourier transform of $\Phi_s(t)$ in Eq. (7), whereas $F(\Delta)$ is the Fourier transform of the relaxation function of the coherence element $\rho_{eg}(t)$. For the common dephasing model corresponding to *exponential decay* of $\rho_{eg}(t)$, $F(\Delta)$ is a Lorentzian spectrum whose width is $\langle\Delta\omega^2\rangle\tau_c$, the product of the mean-square Stark shift and the noisy-field correlation

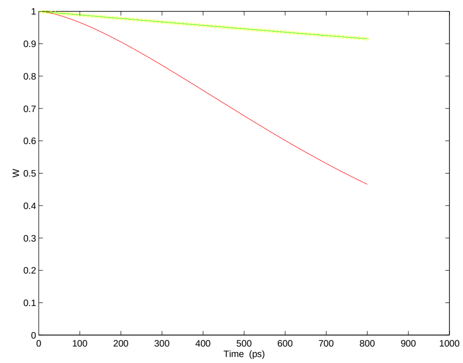

Figure 6: Movie C: Idem, for two-level atom ($\gamma_f = 10^6$ s^{-1}) coupled to waveguide field, with coupling $C^{2/3} = 10^6$ s^{-1} and sharp cutoff. Red curve – uninterrupted evolution at cutoff frequency ($\Delta = 0$); green curve – interrupted evolution at intervals $\tau = 10^{-8}$ s for $\Delta = 0$, yellow dots mark interruption moments.

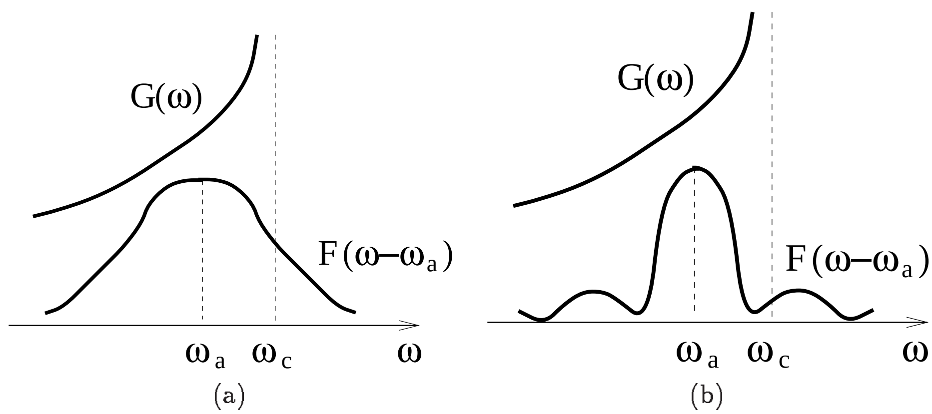

Figure 7: Dependence of effective decay rate κ_s (eq.(14)), on dephasing (relaxation) spectrum $F(\Delta)$ and field reservoir response with cutoff $G(\omega)$: (a) Lorentzian dephasing spectrum; (b) sinc-function spectrum (impulsive measurements).

time. The QZE condition is that *this width be larger than the width of* $G_s(\omega)$ (Fig. 7). The advantage of this realization is that it does not depend on γ_u, and is realizable for any atomic transition. Its importance for molecules is even greater: if we start with a single vibrational level of $|e\rangle$, no additional levels will be populated by this process.

The random ac-Stark shifts described above cause both shifting and broadening of the spectral transition. If we wish to avoid the shifting altogether, we may employ random phase modulation of a driving field that is nearly resonant with the $|e\rangle \leftrightarrow |u\rangle$ transition. If the spectral width of this phase modulation, Γ_ϕ, is much larger than all the other rates specified below, then we can show that

$$\kappa_s \sim 2g_s^2\Gamma_\phi/\Omega^2, \tag{15}$$

where Ω is the driving-field Rabi frequency. When $g_s \ll \Omega \ll \Gamma_\phi$, κ_s is strongly suppressed compared to its zero-field value, and QZE is at its best!

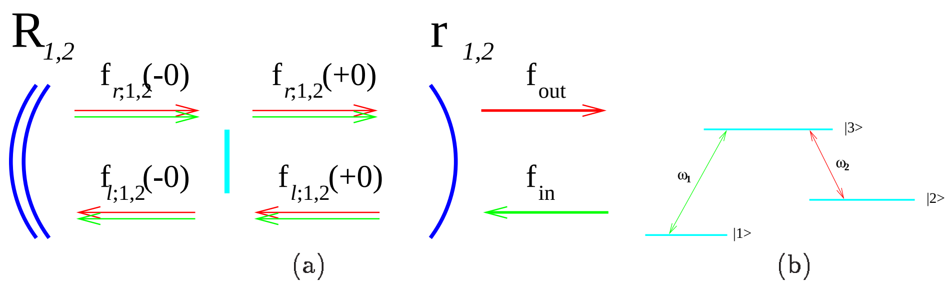

Figure 8: (a) Scheme of location-dependent interference of decay channels in a cavity. Vertical cyan line denotes a thin film, $f_{in(out)}$ denote incoming and outgoing photon, $r(l)$ and 1(2) subscripts denote photon propagation to right (left) near ω_1-green (ω_2-red). (b) Level scheme for atoms in film.

3. State control by location-dependent interference of decay channels

The conventional approaches to state control in a multilevel system are based on interference of absorbed or emitted radiation in two or more channels (transitions), which allows for coherences between field-dressed states of the system[17–25]. In these approaches, spontaneous emission has an adverse effect on interference and its suppression is highly desirable. The adiabatic population transfer method (STIRAP), for example, achieves this goal by preventing the intermediate decaying state from being populated[23]. Here, by contrast, we outline a state-control method[2] wherein spontaneous emission from the intermediate state is, in fact, beneficial, since it exhibits interference between competing channels (transitions), owing to the boundary conditions on the cavity.

We envisage a thin film of atoms or molecules located in a certain position relative to the cavity mirrors. The coupling to the cavity field in the "bad cavity" regime effectively amounts to having a one-dimensional ($1D$) mode continuum, corresponding to spontaneous emission predominantly into one or several Lorentz-broadened modes along the cavity axis. Let us assume that one of the cavity mirrors admits a single photon, a spectrally narrow wavepacket around ω_1. This input condition can be realized by unidirectional mirror transparency near ω_1, or by a measurement verifying that the incoming photon has not been reflected by this mirror. Let us consider the atoms on the thin film to be in the Λ-configuration, such that ω_1 is resonant with the $|1\rangle \leftrightarrow |3\rangle$ transition and ω_2 is resonant with the $|2\rangle \leftrightarrow |3\rangle$ transition (Fig. 8).

Suppose that the task at hand is population transfer from $|1\rangle$ to $|2\rangle$, by means of *resonant absorption* of the ω_1-photon together with spontaneous emission, which is concentrated in the bands around ω_1 and ω_2. The Wigner-Weisskopf solution for this system, at times much longer than the spontaneous lifetime in the cavity γ_c^{-1}, can be obtained upon taking account of the boundary conditions at the mirrors and neglecting the retardation of the photon (due to its travel between the mirrors) relative to the atomic response[2]. This solution satisfies the requirement for complete (100%) population transfer from $|1\rangle$ to $|2\rangle$, provided that

$$\frac{R_1 e^{i\theta_1}(1 + R_2 e^{i\theta_2})(1 + r_2 e^{i\phi_2})}{(1 + R_1 e^{i\theta_1})(1 - r_2 e^{i\phi_2} R_2 e^{i\theta_2})} = \frac{g_1^2}{g_2^2} \tag{16}$$

Here $g_{1(2)}$ are the vacuum field-dipole couplings (vacuum Rabi frequencies) for the $|1\rangle \leftrightarrow |3\rangle$ and $|2\rangle \leftrightarrow |3\rangle$ transitions, respectively; $R_{1(2)}$ and $r_{1(2)}$ are the ω_1- or ω_2-reflectivities of the left-hand and right-hand mirrors, respectively, whereas $\theta_{1(2)}$ ($\phi_{1(2)}$) are the corresponding phase delays accumulated by reflection and round-trip travel to and from the left-hand (right-hand) mirror. Note that r_1, ϕ_1 do not appear in this

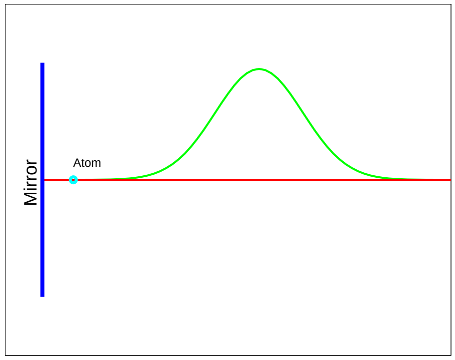

Figure 9: Movie D: Wavepacket conversion from ω_1 input (green envelope) to ω_2 output (red envelope) in cavity under condition (16). $\tau \sim 50\gamma_c^{-1}$ gives an error of $\sim 10^{-4}$ in population transfer. Distance between mirror and atom is assumed to be much less than pulse envelope (shown not to scale), so that the steady state approximation holds for the field between atom and mirror (multiple reflections occur within the pulse propagation time, and are not resolvable on this time scale).

expression, since the complete $|1\rangle \to |2\rangle$ transfer condition must be accompanied by complete conversion of the incident ω_1-photon into the ω_2-photon (Fig. 8).

Movie D traces the evolution of a ω_1-photon wavepacket entering the cavity under condition (16). We can watch the synchronism between its disappearance and the buildup of the exiting ω_2-photon wavepacket.

If we take $R_1 = R_2 = 1$, $\theta_1 = \theta_2 = 2\pi$, i.e., an ideally-reflecting right-hand mirror and an antinode position for the thin film, then we must choose $r_2^2 = |g_1^2 - g_2^2|/(g_1^2 + g_2^2)$ and $\phi_2 = 2\pi$ for $g_1 > g_2$ ($\phi_2 = \pi$ for $g_1 < g_2$). In turn, the g_1/g_2 ratio is determined by the height ratio of the mode Lorentzian around ω_1 and ω_2 and the dipole moments μ_1, μ_2 for these transitions. These conditions impose totally destructive interference between the ω_1-emission and its reflection, and constructive interference between their ω_2-counterparts.

In principle, it is possible to prepare *any superposition* of states $|1\rangle$ and $|2\rangle$ by this method, whereby the intermediate state $|3\rangle$ is excited by a single photon at ω_1 or ω_2, and decays into two channels with appropriately chosen interference between the emitted and reflected amplitudes. Errors in this state control will be incurred by the decay rate γ_b of $|3\rangle$ into the background modes (unconfined by the cavity). In this respect, a cavity with large solid angle confinement, wherein γ_b is partly suppressed, is helpful. In any case, the success of this method depends on the satisfaction of the following condition by the state preparation time (or incident pulse duration) τ:

$$\gamma_b^{-1} \gg \tau \gg \gamma_c^{-1} \tag{17}$$

which implies that the spontaneous emission rate in the cavity, γ_c, must greatly exceed the decay rate γ_b into unconfined modes.

4. Decoherence control by conditionally interfering parallel evolutions

The true challenge of protecting quantum states from decoherence is to devise a simple, low-cost method that would apply *universally* to *any* state, i.e., would not require any prior knowledge of the state. Such a method would be useful for quantum computing and information processing. Here we graphically illustrate the principles of such a method, denoted by us CIPE (Conditionally Interfering Parallel Evolutions), which can control

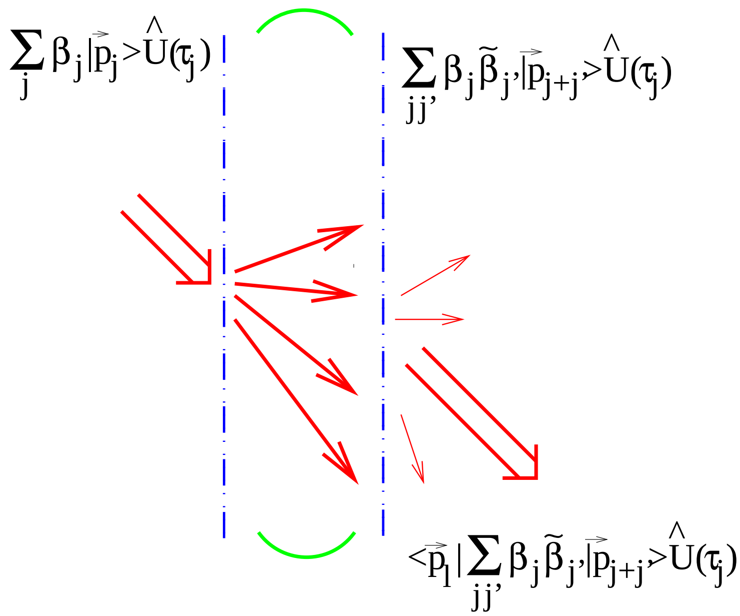

Figure 10: Scheme of CM preparation of superposed parallel field-atom evolutions in a cavity. Red arrows denote atomic momenta. Dash-dotted blue lines denote diffraction gratings.

the evolution of an internal atomic state that is coupled to many nearly *discrete* modes in a cavity.

If the atom is initially in a superposition of its two internal states, and the field in the vacuum state, $|\psi^{(i)}_{A+F}\rangle = (B|e\rangle + C|g\rangle)|0\rangle$, then the atomic interaction with M field modes during a time t results in an *entangled* state

$$U(t)|\psi^{(i)}_{A+F}\rangle \quad = \quad B\left(h(t)|e\rangle|0\rangle + \sum_{m=0}^{M} f_m(t)|g\rangle|1_m\rangle\right) + C|g\rangle|0\rangle, \tag{18}$$

with certain $h(t)$ and $f_m(t)$ dictated by the unitary evolution operator $U(t)$ of the system. The atom-field entanglement in (18) translates, upon tracing over the field degrees of freedom, into *decoherence* of the reduced atomic density matrix. In order to eliminate the decoherence in the atomic state we suggest to *superpose* many *evolution paths* of the atom-field system differing in the *duration* of the atom-field interaction[3]. This can be achieved by conditional measurement (CM) of "ancilla" atomic variables which determine the field-atom interaction time, e.g., the atomic momentum.

By passing the atom through an appropriately designed diffraction grating, before it enters the cavity, we obtain at the exit from the cavity a combined (field-atom) state of the form

$$|\psi^{(f)}_{A+F}\rangle \quad = \quad \sum_j \beta_j|\vec{p}_j\rangle U(t_j)|\psi^{(i)}_{A+F}\rangle. \tag{19}$$

Here the superposed intracavity $A-F$ evolution operators $U(t_j)$ are functions of the evolution time, $t_j = mL/p_{-j}$, which is the corresponding cavity traversal time determined by the momentum projections p_{-j} perpendicular to the cavity axis (Fig. 10), while β_j are the corresponding diffraction amplitudes. The CM projects this state onto an appropriate *superposition* of momentum states. This can be done by passing the atom through a second grating after the cavity whose transmission function is centered at the same momenta $|\vec{p}_j\rangle$, but with different amplitudes $\tilde{\beta}_j$. The CM amounts to detecting the atom in the direction corresponding to $|\vec{p}_l\rangle$. The system obtained by such a CM is then given by

$$|\psi_{A+F}^{(f)}\rangle = \sum_j \alpha_j U(t_j)|\psi_{A+F}^{(i)}\rangle \tag{20}$$

with $\alpha_j = P^{-1/2}\beta_j\beta_{l-j}$, where P is the success probability of the momentum state detection. Tracing over the field degrees of freedom, we find that in order for the final atomic state to be identical with the initial state, namely $\rho_A^{(f)} = \rho_A^{(i)}$, it is *necessary and sufficient* that the following $(M+2)$ conditions be satisfied by the set of α_j's and t_j's parameters:

$$\sum_j \alpha_j = 1 \tag{21}$$

$$\sum_j \alpha_j h(t_j) = 1 \tag{22}$$

$$\sum_j \alpha_j f_m(t_j) = 0 \qquad m = 1..M. \tag{23}$$

Hence, provided that the number of control (momentum-transmission) parameters is sufficiently large, it should be possible in principle to satisfy them simultaneously and thus to ensure that any internal atomic state remains unspoilt in spite of the atom's exposure to many interaction channels. Furthermore, additional control parameters can be used to optimize the CM success probability P. In Fig. (11) we show that the entanglement caused by different evolutions $U(t_j)$, and the ensuing tracing over $M = 20$ field modes yield t_j-dependent decoherence of the initial atomic state $B|e\rangle + C|g\rangle$. By contrast, when this tracing follows the interference of $M+2$ evolution operations $U(t_j)$ with the rightly chosen α_j, it yields an atomic state that is identical to the initial one, i.e., completely protected from decoherence. Considerable deviations from the prescribed conditions on α_j and t_j, i.e., a "failed" measurement, may result in worse decoherence than that associated with each t_j. Hence the need for accurate assignment of the control parameters and for *optimization* of the CM success probability, which is achievable by means of extra control parameters.

5. Conclusions

The present outline and graphical illustration of several methods for combating population decay and state decoherence in cavities underscores their essential traits and allows a comparison of their merits and limitations:
a. The quantum Zeno effect (QZE): Our unified analysis of two-level system coupling to field reservoirs has revealed the general optimal conditions for observing the QZE in various structures (cavities, waveguides, phonon reservoirs, and photonic band structures). We note that the wavefunction collapse notion is not involved here, since the measurement is explicitly described as an act of coherence-breaking. This analysis also clarifies that QZE cannot combat the background-modes contribution to exponential decay. The best way to prevent this contribution is by an AC Stark shift of the resonance well into a photonic band gap. However, if the resonant transition has a considerable width, caused by vibrational-state multiplicity or by inhomogeneous broadening, such that this width exceeds the band-gap width or the AC Stark shift, then we should resort to the QZE strategies outlined above. This method achieves the inhibition of population decay at the expense of increasing decoherence.

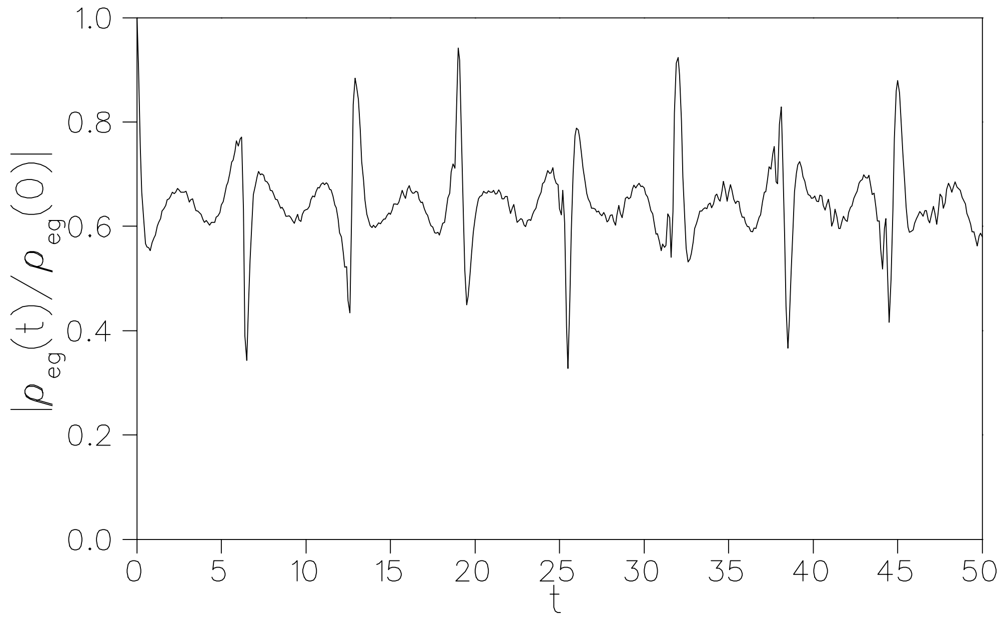

Figure 11: Coherence term $\rho_{eg}(t)$ of the atomic reduced density matrix as a function of evolution time t, showing decoherence of the atomic state with time. By the CIPE method (Eqs.(21)-(23)) we can recreate $\rho_{eg} = 1$ at any desired time t if the required CM is successful.

b. State control by location-dependent interference of decay channels: This method is remarkable in that, rather than try to avoid spontaneous emission by its inhibition (in cavities or in photonic bandgap[16] or by adiabatic off-resonance transfer[23]), it guides this decay into desired channels, and thereby can populate any desired superposition of two (or more) atomic (molecular) states and, correspondingly, of radiation frequencies in the photon wavepacket emerging from the cavity. The limitations of this method are: *(i)* It is appropriate only for thin films, whose location is controllable to within a fraction of a wavelength. *(ii)* As in the case of QZE, it requires the spontaneous emission rates into the axial cavity modes γ_c to be much larger than the corresponding rates into the unconfined "background" modes γ_b. The decay probability of the atom into unconfined modes, $\sim \gamma_b \tau$ during the time τ needed for the creation of the superposition state and the photon wavepacket transformation is the error probability of this state control method. The method is effective only if $1/\gamma_c \ll \tau \ll 1/\gamma_b$.

c. Conditionally interfering parallel evolutions (CIPE): The main merit of this method is that it allows us to control the evolution of $M \gg 1$ coupled degrees of freedom by essentially as many ancilla, which definitely makes it a "low cost" control method. In particular, its applicability to arbitrary atomic superposition states implies that we can use this method as part of the quantum computing/processing chain of operations. By optimizing the CM, we can achieve high success probability, or confidence in the method. Nevertheless, the probabilistic nature of this method remains a limitation. Furthermore, the method does not apply to decay into mode continua, but only to coupling with discrete (albeit large) number of modes.