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Influence of Dissipation On the Stimulated Raman Adiabatic Passage *

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Abstract Using our recently developed quantum dissipation theory, we study the influence of dissipation on the population transfer efficiency between the stimulated Raman adiabatic passage (STIRAP) and the pump-dump passage in a simple three-level Λ system. By comparison, STIRAP is found to be relatively insensitive to the relaxation and fluctuation of the intermediate state. Numerical results also demonstrate the applicability of the new quantum dissipation theory for studying the correlated driving-dissipation dynamics problems.

Key words STIRAP, Quantum dissipation

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1 Introduction

Stimulated Raman adiabatic passage (STIRAP) is an efficient and robust technique for population transfer between states and has been extensively studied both experimentally and theoretically for years[1]. STIRAP that involves no real occupation in the intermediate state(s) is characterized by its counterintuitive field configuration in which the Stokes (dump) pulse proceeds the pump pulse. The STIRAP scheme was first put forward theoretically by Hioe and Eberly and co-workers[2] and later realized experimentally by Bergmann and coworkers[3]. Applications have also been made in some atomic cooling[4] and Bose-Einstein condensate[5] systems. Recently, some related interests have been extended to multi-level systems[6], continuum systems[7], and selective population transfer for degenerate states[8].

Despite its success in various systems, the theoretical studies of the STIRAP control scheme are by far largely only for gas phase systems. The main obstacle here is due to the lack of a reliable formulation to treat the correlated driving and non-Markovian dissipation. In this work, we shall exploit our recently developed quantum dissipation theory[9] to investigate the influence of dissipation on STIRAP.

2 Background of quantum dissipation theory

The object of study in quantum statistical mechanism is the reduced density operator $\rho(t) \equiv tr_B[\rho_T(t)]$, defined as the partial trace of the total system-plus-bath density operator $\rho_T(t)$ over the bath degrees of freedom. The total Hamiltonian assumes the form of $H_T = H(t) + QF(t)$. Here, $H(t) = H_s + \mu\varepsilon(t)$, with H_s being the system Hamiltonian, μ the dipole operator, and $\varepsilon(t)$ the laser field. Q is a system operator, while $F(t)$ denotes the generalized Langevin force which is a stochastic operator in the bath subspace. We have recently construct a so-called complete second-order quantum dissipation theory (CS-QDT) in which the system-bath interaction is treated rigorously at second-order cumulant level for

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both reduced dynamics and initial canonical boundary condition. The final CS-QDT reads as (setting $\hbar \equiv 1$)[9\]

$$\dot{\rho}(t) = -i[H(t), \rho(t)] - [Q, \tilde{Q}(t)\rho(t) - \rho(t)\tilde{Q}^{\dagger}(t)] \quad (1a)$$

$$\tilde{Q}(t) = \int_{-\infty}^{t} d\tau C(t-\tau) G(t,\tau) Q G^{\dagger}(t,\tau) \quad (1b)$$

Here, $C(t)$ is the force-force correlation function[9,10], describing the effects of bath interaction. It satisfies the detailed-balance relation, $C^*(t) = C(t - i\beta)$, which in the frequency-domain reads as $\hat{C}(\omega) = \exp(\beta\omega)\hat{C}(-\omega)$. Here, $\beta = 1/(k_B T)$, with k_B being the Boltzmann constant and T the temperature.

Note that Eq.(1) is of a coupled differential-integral form whose numerical implementation is rather expensive. To facilitate this problem, we introduce the spectral density function,

$$J(\omega) \equiv \frac{1}{2}[\hat{C}(\omega) - \hat{C}(-\omega)]$$

and adopt the Meier-Tannor's parameterization algorithm[9,11]:

$$J(\omega) = \sum_k \frac{p_k \omega}{|\omega^2 - (\Omega_k + i\Gamma_k)^2|^2} \equiv \sum_k \frac{p_k \omega}{|\omega^2 - z_k^2|^2} \quad (2)$$

p_k is a real, Ω_k and Γ_k are positive fitting parameters. The bath correlation function can now be expressed as[9,11]:

$$\begin{aligned} C(t \geqslant 0) &= \frac{1}{\pi}\int_{-\infty}^{\infty} d\omega \frac{J(\omega)\exp(-i\omega t)}{1 - \exp(-\beta\omega)} \\ &= \sum_k \frac{p_k}{4\Omega_k \Gamma_k}\left[\frac{\exp(iz_k t)}{\exp(\beta z_k) - 1} + \frac{\exp(-iz_k^* t)}{1 - \exp(-\beta z_k^*)}\right] - \frac{1}{\pi}\sum_{n=1}^{\infty}\frac{1}{n}\tilde{J}(\tilde{\omega}_n)\exp(-\tilde{\omega}_n t) \\ &\equiv \sum_m \nu_m \exp(-\zeta_m t) \end{aligned} \quad (3)$$

Here, $\tilde{\omega}_n \equiv 2\pi n/\beta$ and $\tilde{J}(\omega) \equiv i\omega J(-i\omega)$. For finite temperatures, the summation over $\tilde{\omega}_n$ can be truncated at some value N'. As results, we can recast Eq.(1b) as:

$$\tilde{Q}(t) \equiv \sum_m K_m(t) \quad (4a)$$

$$\dot{K}_m(t) = \nu_m Q - \zeta_m K_m(t) - i[H(t), K_m(t)] \quad (4b)$$

Thus, Eqs.(1a) and (4) constitute a closed set of equations of motion which can be numerically propagated efficiently. The initial conditions to these equations, which relate to the thermal equilibrium reduced density operator, can be obtained by evaluating the stationary solutions in the absence of external fields[9].

3 Numerical results and discussion

Shown in Fig.1 is a three-level Λ system typically used in the STIRAP study. The levels $|1\rangle$ and $|2\rangle$ are coupled by the pump field ε_p, while $|2\rangle$ and $|3\rangle$ are coupled by the Stokes (dump) field ε_s. There is no direct coupling between $|1\rangle$ and $|3\rangle$. The carrier frequency of the pump (Stokes) field is set to be at resonance with the $2 \leftarrow 1$ ($3 \leftarrow 2$) transition frequency ω_p ($\omega_s = 0.2\omega_p$), and $k_B T = 0.16\omega_p$. The pump-dump fields are Gaussian pulses, with a same Gaussian width of $1000/\omega_p$ and a same dipole interaction strength of
$0.02\omega_p$ for STIRAP or $0.0013\omega_p$ for the conventional pump-dump scheme. The Stokes/dump delay time t_d is set at $\pm 2000/\omega_p$. In the STIRAP scheme, $t_d < 0$ and the Stokes field is prior to the pump field, which is in contrast with the conventional pump-dump scheme where $t_d > 0$.

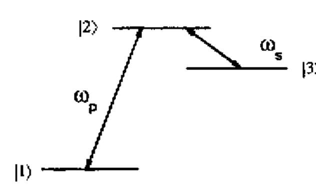

Fig.1 The typical three-level Λ system in a STIRAP control scheme (see the text for details)

Fig. 2 depicts the transient populations of the three levels in both the STIRAP (left panels) and the conventional pump-dump (right panels) control schemes. The top panels are for the dissipation-free gas phase system. The others are for three representative cases of dissipation: two involve the T_1-relaxation between the pumping levels (γ_1^p) and between the Stokes dumping levels (γ_1^s), respectively, and one involves a pure T_2-dephasing (γ_2). These three cases of dissipation can be described with Q_p, Q_s and Q_2, respectively, for the operator Q in Eqs. (1a) and (4). Their nonzero matrix elements are only $(Q_p)_{12} = (Q_p)_{21} = (Q_s)_{23} = (Q_s)_{32} = (Q_2)_{22} = 1$. The interaction bath spectral density $J(\omega)$ in each case of dissipation assumes to be Ohmic, $J(\omega) = \eta\omega\exp(-|\omega|/\omega_c)$, with the cut-off frequency $\omega_c = \omega_p$. The system-bath coupling strength η is chosen by setting that $\gamma_1^p = C(\omega_p) = 0.001\omega_p$, $\gamma_1^s = C(\omega_s) = 0.001\omega_p$, or $\gamma_2 = C(0) = 0.001\omega_p$, respectively, for each of the three cases of dissipation. The parameters for fitting the Ohmic $J(\omega)$ in the form of Eq. (3) are taken from Ref. [11].

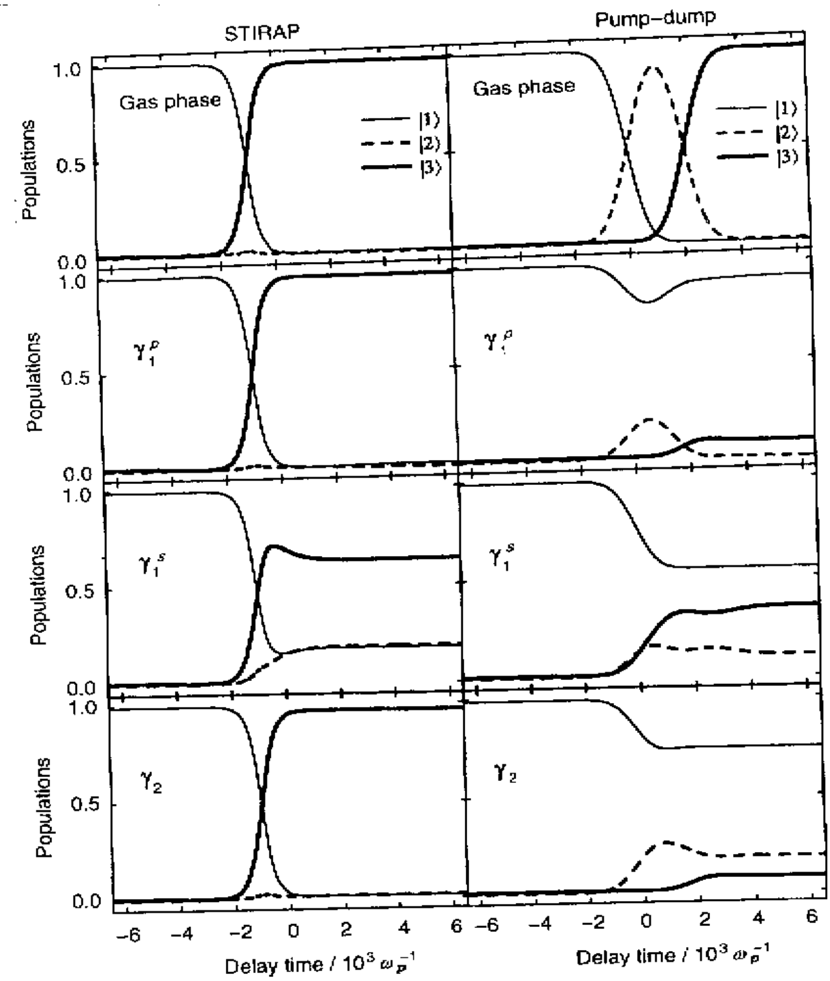

Fig. 2 Population transfer in comparison between the STIRAP and the conventional pump-dump control schemes Panels from top to bottom are for the dissipation-free gas phase, the T_1-relaxation between the $1\leftrightarrow 2$ pumping levels (γ_1^p), the T_1-relaxation between the $2\leftrightarrow 3$ Stokes dumping levels (γ_1^s), and the pure T_2-dephasing (γ_2) systems, respectively (see the text for details).

It can be seen from the left panels of Fig. 2 that the γ_1^p-relaxation and γ_2-dephasing affect little on the quality of STIRAP. The fact that the γ_1^s-relaxation decreases the transfer efficiency can be explained by the establishment of the thermal equilibrium between the Stokes levels, i. e., the nonradiative decay between $|2\rangle$ and $|3\rangle$. The STIRAP processes in all these left four panels do not radiatively promote the population in the intermediate level $|2\rangle$. In comparison, the conventional pump-dump processes (the right panels of Fig. 2) not only radiatively pump the population onto the intermediate level, but also are very sensitive to the fluctuating environment, which deteriorates dramatically the quality of the desired 3←1 transfer.

4 Summary

We have demonstrated the robustness of the STIRAP control scheme against both the T_1-relaxation and T_2-dephasing. Note that the STIRAP scheme is characterized by the configuration in which the prior Stokes pulse should also overlap with the pump pulse and both fields are in the strong response regime[1-3]. The STIRAP mechanism[1-3] that supports a robust population transfer is still valid in the presence of dissipation. To treat this type of correlated driving-dissipation dynamics problems, the traditional quantum dissipation theory such as Bloch-Redfield quantum master equation is found numerically diverged due to the approximations that involved[12]. The CS-QDT outlined in Sec. 2 is however proven to be very stable, leading to physically meaningful results here. Implementation of the new quantum dissipation theory to other types of optical control schemes that have only been considered for gas phases by far will be made in future.

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耗散对受激拉曼绝热转移过程的影响*

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摘 要: 应用最近发展的量子耗散理论,研究了耗散对简单三能级体系的受激拉曼绝热转移过程的影响,并与 pump-dump 过程比较.计算结果表明,受激拉曼绝热转移的机制能很好地抑制中间态的弛豫与涨落的影响.数值结果也表明了新的量子耗散理论可以正确地描述场与耗散相互耦合的动力学问题.

关键词: 受激拉曼绝热转移过程;量子耗散

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