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Plasmon-assisted quantum control of distant emitters

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ABSTRACT

We show how to generate and control the correlations in a set of two distant quantum emitters coupled to a one-dimensional dissipative plasmonic waveguide. An external laser field enhances the dimer's steady-state correlations and allows an active control (switching on/off) of nonclassical correlations. The plasmon-assisted dipolar-interacting qubits exhibit persistent correlations, which in turn can be decoupled and made to evolve independently from each other. The setup enables long-distance (~ 1 μm) qubit control that works for both resonant and detuned emitters. For suitable emitter initialization, we also show that the quantum correlation is always greater than the classical one.

1. Introduction

The construction of large quantum networks with controllable long-distance coupling of qubits is one of the key goals in quantum information science [1]. To generate the required qubit–qubit correlations (microwave or optical) photons are typically used. Recently, a novel approach based on single quantum emitters coupled to one-dimensional plasmonic waveguides has been proposed, and entanglement of qubits mediated by surface plasmons in such a setup has been explored [2,3]. Surface plasmons are collective excitations that have become an important physical resource in many applications in physics, chemistry, and materials science [4,5]. In particular, the coupling of single emitters to plasmonic structures [6–8] has attracted substantial attention, because it allows the manipulation of the emission properties as well as the enhancement of the interaction between quantum emitters in the vicinity of these structures [9–11].

Non-local quantum correlations and entanglement in quantum systems under the action of decoherence effects have been arduously investigated not only because of their fundamental physical implications (see e.g., [12–14]), but also because of their utmost relevance to the development of novel quantum technologies [1]. The influence of the system-bath coupling over the reduced system's quantum correlations has been studied in different physical systems such as quantum dots [15], superconducting qubits [16], atoms and photons [17], and biomolecular systems [18], to cite just but a few. We have recently reported on the distribution of classical and quantum correlations that arise in a bipartite emitter system coupled to a plasmonic waveguide [19]. We have shown that such quantum correlations are more robust through the dissipative dynamics than the classical ones, for several experimentally accessible scenarios, and that the emitters collective properties that arise from the interaction with the plasmons allow an additional degree of quantum control on the correlations dynamics [19,20].

In this Letter, we give a protocol for actively enabling a plasmon-assisted long-distance (about 1 μm) qubit conditional dynamics between emitters which are externally-driven by a coherent laser field. Such a quantum mechanism works for both resonant and detuned qubits, and is optimized by laser pumping and by tailoring the separation between the emitters. Moreover, from the dimer's conditional entropy, we analytically identify specific conditions for which the quantum correlation is always greater than the classical correlation (cf. [21–23]).

In Section 2 we introduce the quantifiers of quantum and classical correlations, and the corresponding quantum master equation used to describe the emitters' dissipative dynamics is given in Section 3. The quantum control that arises from the correlations

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dynamics is set for: resonant emitters without laser excitation (Section 4), resonant emitters under the action of a coherent laser field (Section 5), and for emitters with different transition energies (Section 6). Concluding remarks are given in Section 7.

2. Quantum correlations

The quantum mutual information $\mathcal{I}(\rho_{AB}) = S(\rho_A) + S(\rho_B) - S(\rho_{AB})$ gives a measure of the total correlation in a bipartite qubit state ρ_{AB} [23–25], and this may be separated into purely quantum–$\mathcal{Q}(\rho_{AB})$ (e.g., via the quantum discord $\mathcal{D}(\rho_{AB})$) [25–27] and classical correlation–$\mathcal{C}(\rho_{AB})$ [21]: $\mathcal{I}(\rho_{AB}) = \mathcal{Q}(\rho_{AB}) + \mathcal{C}(\rho_{AB})$. The details of calculation of $\mathcal{C}$ and $\mathcal{Q}$ are left to Appendix A. The von Neumann entropy $S(\rho) = -\operatorname{Tr}\rho\log_2\rho$, and $\rho_{A(B)} = \operatorname{Tr}_{B(A)}\rho_{AB}$ is the reduced density operator of the partition $A(B)$. Discord has been linked to the computational speedup in an efficient model of quantum computation [28] and has been pinpointed as a valuable resource in quantum information protocols [29,28].

Although Lindblad conjectured that, for any quantum state, $\mathcal{C}(\rho) \geq \mathcal{Q}(\rho)$ [23,21,22], we give an entropy condition for which the quantum correlation is always greater than the classical one [30,20]: since $\mathcal{I} = \mathcal{Q} + \mathcal{C}$, we ask if the inequality $\mathcal{I} - 2\mathcal{C} \geq 0$ is ever met. The sought entropy bound reads

$$2\mathcal{S}(\rho_{A|\Pi_j^B}) - S(\rho_{AB}) + S(\rho_B) - S(\rho_A) \geq 0, \tag{1}$$

where $\mathcal{S}(\rho_{A|\Pi_j^B}) = \min_{\{\Pi_j^B\}}\{\sum_j p_j S(\rho_{A|\Pi_j^B})\}$, $S(\rho_{A|\Pi_j^B})$ is the entropy associated to the density matrix of subsystem A after the measure. Eq. (1) is indeed satisfied, at all times, for suitable resonant emitters.

To distinguish the quantum correlations that arise for entangled states and separable states, we quantify the emitters' entanglement via the entanglement of formation $\mathcal{E_F}$. For two-qubit systems, $\mathcal{E_F}(\rho_{AB}) = \varepsilon([1+\sqrt{1-C^2(\rho_{AB})}]/2)$, where $\varepsilon(r) = -r\log_2 r - (1-r)\log_2(1-r)$ denotes the binary entropy function, and the concurrence $C(\rho_{AB}) = \max\{0, \lambda_1 - \lambda_2 - \lambda_3 - \lambda_4\}$, where the λ_i's are, in decreasing order, the eigenvalues of the matrix $\sqrt{\rho_{AB}\tilde{\rho}_{AB}}$; $\tilde{\rho}_{AB} = (\sigma_y \otimes \sigma_y)\bar{\rho}_{AB}(\sigma_y \otimes \sigma_y)$, $\bar{\rho}_{AB}$ is the elementwise complex conjugate of ρ, and σ_y is the Pauli matrix [31]. This entanglement metric is of entropic character and can be compared on the same grounds with the discord [19].

3. Plasmon–emitter master equation

We consider a pair of distant emitters that act as a dipole–dipole two-qubit system via the interaction with the plasmon modes in a metallic nanostructure. The total Hamiltonian of the system plus environment can be written as ($\hbar$ is the reduced Planck's constant):

$$\hat{H} = \hat{H}_S + \hat{H}_E + \hat{H}_{int} = \sum_{i=1}^{2}\frac{1}{2}\hbar\omega_i\hat{\sigma}_z^{(i)} + \int d^3\mathbf{r}\int_0^\infty d\omega\hbar\omega\hat{f}^\dagger_\omega(\mathbf{r})\hat{f}_\omega(\mathbf{r}) - \sum_{i=1}^{2}\hat{\mu}_i\cdot\hat{E}(\mathbf{r}_i), \tag{2}$$

where the first two terms denote the free energy of the two-qubit system and the free energy of the electromagnetic field represented by a bosonic bath, respectively; $\hat{\sigma}_z^{(i)}$ is the z Pauli matrix associated to the i-th emitter of transition frequency ω_i, and $\hat{f}_\omega(\mathbf{r})$, and $\hat{f}^\dagger_\omega(\mathbf{r})$ are the bosonic excitation operators of the quantized electromagnetic field with the usual commutation relations $[\hat{f}_\omega(\mathbf{r}), \hat{f}^\dagger_{\omega'}(\mathbf{r}')] = \delta(\omega-\omega')\delta(\mathbf{r}-\mathbf{r}')$, and $[\hat{f}_\omega(\mathbf{r}), \hat{f}_{\omega'}(\mathbf{r}')] = [\hat{f}^\dagger_\omega(\mathbf{r}), \hat{f}^\dagger_{\omega'}(\mathbf{r}')] = 0$. The last term represents the interaction between the dipole operator $\hat{\mu} = \boldsymbol{\mu_i}\hat{\sigma}_+^{(i)} + \boldsymbol{\mu_i^*}\hat{\sigma}_-^{(i)}$ and the quantized field operator $\hat{E}(\mathbf{r}) = \hat{E}^+(\mathbf{r}) + \text{H.c.}$, where $\sigma_+^{(i)} = |1_i\rangle\langle 0_i|$ ($\sigma_-^{(i)} = |0_i\rangle\langle 1_i|$) are the raising (lowering) Pauli operators acting on emitter i ($|0\rangle$ and $|1\rangle$ denote the ground and first excited state which represent the qubit computational basis), and $\hat{E}^+(\mathbf{r}) = \int_0^\infty d\omega\hat{E}(\mathbf{r},\omega)$ is the positive frequency part,

$$\hat{E}(\mathbf{r},\omega) = \mathrm{i}\sqrt{\frac{\hbar}{\pi\epsilon_0}}\frac{\omega^2}{c^2}\int d^3\mathbf{r}'\sqrt{\epsilon''(\mathbf{r}',\omega)}\boldsymbol{G}(\mathbf{r},\mathbf{r}',\omega)\hat{f}_\omega(\mathbf{r}').$$

The Green's tensor $\boldsymbol{G}(\mathbf{r},\mathbf{r}',\omega)$ satisfies the Maxwell–Helmholtz wave equation, supports the electromagnetic interaction from $\mathbf{r}'$ to $\mathbf{r}$, and contains all the information about the coherent and incoherent properties of the system, i.e., about the dipole–dipole shift and the different channels of radiation through the vacuum (far field) and through the metal. $\epsilon''(\mathbf{r},\omega)$ is the imaginary part of the electric permittivity of the metal, and this is considered as a constant value corresponding to the permittivity of the silver at the operational wavelength described here, and ϵ_0 is the permittivity of the vacuum.

The dynamics of the total system (emitters plus electromagnetic field) can be derived in the Schrödinger picture from the Liouville–von Neumann equation $\dot{\rho}_{S-E} = -\frac{\mathrm{i}}{\hbar}[\hat{H}, \rho_{S-E}]$, which, in the interaction picture (denoted by the superscript I), can be written as the following integro-differential equation

$$\dot{\rho}^I_{S-E}(t) = -\frac{\mathrm{i}}{\hbar}\big[\hat{H}^I_{int}(t), \rho^I_{S-E}(0)\big] - \frac{1}{\hbar^2}\int_0^t dt'\big[\hat{H}^I_{int}(t), \big[\hat{H}^I_{int}(t'), \rho^I_{S-E}(t')\big]\big], \tag{3}$$

where $\rho^I_{S-E}(t) = e^{\mathrm{i}(\hat{H}_S+\hat{H}_E)t/\hbar}\rho_{S-E}e^{-\mathrm{i}(\hat{H}_S+\hat{H}_E)t/\hbar}$, and $\hat{H}^I_{int}(t) = e^{\mathrm{i}(\hat{H}_S+\hat{H}_E)t/\hbar}\hat{H}_{int}e^{-\mathrm{i}(\hat{H}_S+\hat{H}_E)t/\hbar}$. Here, we consider atom-like (small) emitters with an operational wavelength in the optical frequency regime (we use $\lambda_0 = 640$ nm) coupled to a broadband plasmonic waveguide. We perform the Born–Markov approximation ($\rho^I_{S-E}(t') = \rho^I(t')\otimes\rho_E(0)$, and $\rho^I(t') \to \rho^I(t)$) in order to solve Eq. (3) due to the weak coupling between the emitters and the plasmonic electromagnetic field [2,3,11,32]. Additionally, we avoid rapidly oscillating terms much higher than ω_i by applying the rotating wave approximation (RWA) to the interaction part of the Hamiltonian in Eq. (2), such that $\hat{H}_{int} = -\boldsymbol{\mu}_1\sigma_+^{(1)}\hat{E}^\dagger(\mathbf{r}_1) - \boldsymbol{\mu}_2\sigma_+^{(2)}\hat{E}^\dagger(\mathbf{r}_2) + \text{H.c}$.

Tracing out Eq. (3) over the environment degrees of freedom, and going back to the Schrödinger picture we can, under the above assumptions, describe the dimer's dissipative dynamics by means of the following quantum master equation [11,20,33]

$$\dot{\rho} = \frac{\mathrm{i}}{\hbar}[\rho, \hat{H}_{\text{eff}}] - \sum_{i,j=1}^{2}\frac{\Gamma_{ij}}{2}\big(\rho\sigma_+^{(i)}\sigma_-^{(j)} + \sigma_+^{(i)}\sigma_-^{(j)}\rho - 2\sigma_-^{(i)}\rho\sigma_+^{(j)}\big), \tag{4}$$

where the effective dimer's Hamiltonian $\hat{H}_{\text{eff}} = \hat{H}_S + \hat{H}_{12} + \hat{H}_L$ contains the coherent dipole–dipole shift $\hat{H}_{12} = \frac{1}{2}\hbar V(\sigma_x^{(1)}\otimes\sigma_x^{(2)} + \sigma_y^{(1)}\otimes\sigma_y^{(2)})$, and the laser–qubit interaction $\hat{H}_L$ (see Section 5). The strength of the effective dipole–dipole interaction is calculated as:

$$V = \frac{1}{\pi\epsilon_0 c^2\hbar}\mathcal{P}\int_0^\infty d\omega\frac{\omega^2\,\mathrm{Im}[\boldsymbol{\mu}_1^*\boldsymbol{G}(\omega,\mathbf{r}_1,\mathbf{r}_2)\boldsymbol{\mu}_2]}{\omega-\omega_0}, \tag{5}$$