

Coherent and incoherent dynamics in excitonic energy transfer: Correlated fluctuations and off-resonance effects

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We study the nature of the energy transfer process within a pair of coupled two-level systems (donor and acceptor) subject to interactions with the surrounding environment. Going beyond a standard weak-coupling approach, we derive a master equation within the polaron representation that allows for the investigation of both weak and strong system-bath couplings, as well as reliable interpolation between these two limits. With this theory, we are then able to explore both coherent and incoherent regimes of energy transfer within the donor-acceptor pair. We elucidate how the degree of correlation in the donor and acceptor fluctuations, the donor-acceptor energy mismatch, and the range of the environment frequency distribution impact upon the energy transfer dynamics. In the resonant case (no energy mismatch) we describe in detail how a crossover from coherent to incoherent transfer dynamics occurs with increasing temperature [A. Nazir, *Phys. Rev. Lett.* **103**, 146404 (2009)], and we also explore how fluctuation correlations are able to protect coherence in the energy transfer process. We show that a strict crossover criterion is harder to define when off-resonance, though we find qualitatively similar population dynamics to the resonant case with increasing temperature, while the amplitude of coherent population oscillations also becomes suppressed with growing site energy mismatch.

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I. INTRODUCTION

A fascinating series of recent experiments demonstrating signatures of quantum coherence in the energy transfer dynamics of a variety of systems^{1–9} has sparked renewed interest in modeling excitation energy transfer beyond standard methods.^{10–19} This process, which occurs when energy absorbed at one site (the donor) is transferred to another nearby site (the acceptor) via a virtual photon,²⁰ is often considered to be incoherent; the result of weak donor-acceptor interactions, treated perturbatively using Fermi's golden rule.^{21,22} However, though this approach has proved to be immensely successful when applied in many situations,^{23,24} accounting for *quantum coherence* within the energy transfer dynamics requires an analysis beyond straightforward perturbation theory in the donor-acceptor interaction.

An alternative starting point for investigations into coherent energy transfer is to treat the system-environment interaction as a perturbation instead. Such weak-coupling theories, often referred to as being of Redfield or Lindblad type depending upon the approximations made in their derivation,²⁵ have been successfully applied to elucidate a number of effects that could be at play in multisite donor-acceptor complexes. Examples include studying the interplay of coherent dynamics and dephasing in promoting efficient energy transfer in quantum aggregates,^{26–32} exploring the role of environmental correlations in tuning the energy transfer process,³³ and extensions to assess the potential importance of non-Markovian effects.^{34,35}

Nevertheless, to properly understand the transition from coherent to incoherent energy transfer which occurs as the system-environment coupling or temperature is increased,^{11,36–38} it is necessary to be able to describe the system dynamics beyond either of these limiting cases.^{24,39,40} Building on earlier work,^{36,41–43} a number of methods have

been put forward to accomplish this. For example, modifications to both Redfield^{44–47} and Förster^{48–50} theories have extended the range of validity of both approaches. Moreover, it is possible to define an alternative perturbation term through the small polaron transformation,⁵¹ which under certain conditions allows interpolation between the Redfield and Förster limits.^{11–13} For particular forms of system-environment interaction, this can also be achieved through the hierarchical equations of motion technique.^{16,52} Numerically exact calculations, based, for example, on path integral,^{15,53} numerical renormalization group⁵⁴ and density matrix renormalization group¹⁴ methods, have also been applied to study energy transfer beyond perturbative approaches.

In this work, we investigate the conditions under which coherent or incoherent motion is expected to dominate the energy transfer dynamics of a model donor-acceptor pair. Following Ref. 11, we employ a Markovian master equation derived within the polaron representation for this purpose since it allows for a consistent analysis of the dynamics from weak to strong system-bath coupling (or, equivalently, low to high temperatures).^{51,55} In addition to presenting a full derivation of the theory, we also extend it to explore, in detail, the important effects of donor-acceptor energy mismatch, deriving analytical forms for the dissipative dynamics valid over a large range of parameter space. Furthermore, we move beyond the scaling limit studied in Ref. 11 to consider an environment frequency distribution of finite extent, characterized by a high-frequency cutoff in the bath spectral density. In the resonant case (no energy mismatch) we define a strict crossover temperature above which the energy transfer dynamics ceases to be coherent.¹¹

Of particular practical interest is the role played by correlations between the donor and acceptor environmental fluctuations, suggested as one mechanism by which quantum

coherence may survive in the energy transfer process under otherwise adverse conditions^{1,4,9,11,53,56–60} (though see, for example, Refs. 61 and 62 for alternatives). These correlations are also easily treated within our formalism, through position-dependent couplings between the system and the common environment. As the donor and acceptor are brought closer together, there comes a point at which their separation becomes comparable to, or smaller than, the wavelength of relevant modes in the bath. As this happens, fluctuations at each site become ever more correlated, and dephasing effects are suppressed. We shall show, consequently, that as the level of correlation increases, so too does the crossover temperature to the incoherent regime. Hence, strong correlations lead to the survival of coherence at high temperatures.

Off-resonance, we find that it is less straightforward to define a crossover temperature. In contrast to the resonant case, for sufficient energy mismatch between the donor and acceptor, increasing the temperature causes the amplitude of the coherent contribution to decrease, though not to disappear altogether. In principle, it then becomes possible for a coherent component to exist in the dynamics at all but infinite temperatures. Although we are then unable to define a crossover in quite the same way, we still find that bath correlations have a qualitatively similar effect to the resonant case, protecting coherence in the transfer process.

The paper is organized as follows. In Sec. II we introduce our model, and derive a master equation describing the donor-acceptor dynamics within the polaron representation. Section III considers the resonant case and the coherent-incoherent crossover. In Sec. IV we investigate off-resonant energy transfer and obtain analytic expressions for the dynamics in a number of limits. Finally, in Sec. V we summarize our results.

II. POLARON TRANSFORM MASTER EQUATION

A. The system and polaron transformation

We consider a donor-acceptor pair ($j = 1, 2$), each site of which is modeled as a two-level system with ground state $|G\rangle_j$, excited state $|X\rangle_j$, and energy splitting ϵ_j . The pair interact via Coulombic energy transfer with strength V , which is responsible for the transfer of excitation from one site to the other. We label the state corresponding to a single excitation on site 1 as $|1\rangle \equiv |XG\rangle$, and that on site 2 as $|2\rangle \equiv |GX\rangle$. The environment surrounding the donor-acceptor pair is modeled as a common bath of harmonic oscillators, coupled linearly to the excited state of each site. The total system-bath Hamiltonian in the single excitation subspace (in which energy transfer occurs) is therefore written (where $\hbar = 1$)

$$H_{\text{SUB}} = \epsilon_1 |1\rangle \langle 1| + \epsilon_2 |2\rangle \langle 2| + V(|1\rangle \langle 2| + |2\rangle \langle 1|) + |1\rangle \langle 1| B_z^{(1)} + |2\rangle \langle 2| B_z^{(2)} + \sum_{\mathbf{k}} \omega_{\mathbf{k}} b_{\mathbf{k}}^\dagger b_{\mathbf{k}}, \quad (1)$$

where the bath is described by creation (annihilation) operators $b_{\mathbf{k}}^\dagger$ ($b_{\mathbf{k}}$) with corresponding angular frequency $\omega_{\mathbf{k}}$, and wave vector \mathbf{k} . The bath operators are given by $B_z^{(j)} = \sum_{\mathbf{k}} (g_{\mathbf{k}}^{(j)} b_{\mathbf{k}}^\dagger + g_{\mathbf{k}}^{(j)*} b_{\mathbf{k}})$, with coupling constants $g_{\mathbf{k}}^{(j)}$. As in Ref. 11, we shall consider the case in which each site is coupled to the bosonic

bath with the same magnitude $|g_{\mathbf{k}}|$, but make the separation between the sites explicit through position-dependent phases in the coupling constants of the form $g_{\mathbf{k}}^{(j)} = |g_{\mathbf{k}}| e^{i\mathbf{k} \cdot \mathbf{r}_j}$, with \mathbf{r}_j being the position of site j . As we shall see, this form of coupling gives rise to correlations between the bath influences experienced at each site, allowing a range of totally correlated, partially correlated, and completely uncorrelated fluctuations to be explored.^{11,53}

A standard weak-coupling approach to the system dynamics would now be to derive a master equation for the evolution of the reduced system density operator under the assumption that the system-bath interaction terms, as written in Eq. (1), can be treated as weak perturbations.^{25,63} In this work, we shall instead derive a master equation describing the donor-acceptor energy transfer dynamics in the (now widely used) polaron representation,^{11–13,43,51,55,64} whereby we displace the bath oscillators depending on the system state. We may then identify alternative perturbation terms, which can be small over a much larger range of parameter space than those in the original representation. In particular, the polaron framework allows us to reliably explore from weak (single-phonon) to strong (multiphonon) coupling regimes between the system and the bath, provided that the energy transfer interaction V does not become the largest energy scale in the problem (in which case the full polaron displacement is no longer appropriate⁶⁵), and that there is no infrared divergence in its bath-renormalized value V_R (see Eq. (7) below).⁶⁶ In contrast, with a weak system-bath coupling treatment we would only be able to probe single-phonon bath-induced processes, and hence not be able to properly explore the crossover from coherent to incoherent dynamics in which we are primarily interested.

To proceed, we thus apply a unitary transformation which displaces the bath oscillators according to the location of the excitation. Defining $H_P = e^S H_{\text{SUB}} e^{-S}$, where

$$S = |1\rangle \langle 1| P(g_{\mathbf{k}}^{(1)}/\omega_{\mathbf{k}}) + |2\rangle \langle 2| P(g_{\mathbf{k}}^{(2)}/\omega_{\mathbf{k}}), \quad (2)$$

with bath operators $P(\alpha_{\mathbf{k}}) = \sum_{\mathbf{k}} (\alpha_{\mathbf{k}} b_{\mathbf{k}}^\dagger - \alpha_{\mathbf{k}}^* b_{\mathbf{k}})$, results in the polaron transformed spin-boson Hamiltonian⁵¹ $H_P = H_0 + H_I$, with

$$H_0 = \frac{\epsilon}{2} \sigma_z + V_R \sigma_x + \sum_{\mathbf{k}} \omega_{\mathbf{k}} b_{\mathbf{k}}^\dagger b_{\mathbf{k}}, \quad (3)$$

and

$$H_I = V(B_x \sigma_x + B_y \sigma_y). \quad (4)$$

Here, the bias $\epsilon = \epsilon_1 - \epsilon_2$, gives the energy difference between the donor and acceptor, while the Pauli operators are defined in a basis in which $\sigma_z = |1\rangle \langle 1| - |2\rangle \langle 2| = |XG\rangle \langle XG| - |GX\rangle \langle GX|$. The bath operators appearing in Eq. (3) are constructed as $B_x = (1/2)(B_+ + B_- - 2B)$ and $B_y = (i/2)(B_+ - B_-)$, where

$$B_{\pm} = \prod_{\mathbf{k}} D\left(\pm \frac{(g_{\mathbf{k}}^{(1)} - g_{\mathbf{k}}^{(2)})}{\omega_{\mathbf{k}}}\right), \quad (5)$$

with displacement operators $D(\pm\alpha_{\mathbf{k}}) = \exp[\pm(\alpha_{\mathbf{k}} b_{\mathbf{k}}^\dagger - \alpha_{\mathbf{k}}^* b_{\mathbf{k}})]$. Note that the interaction terms in Eq. (4) therefore depend upon the difference in donor and acceptor system-bath couplings $g_{\mathbf{k}}^{(1)}$ and $g_{\mathbf{k}}^{(2)}$, respectively. Importantly, the term driving coherent energy transfer in Eq. (3) will not be treated

perturbatively, though it does now have a bath-renormalized strength, $V_R = BV$, where

$$B = \exp \left[- \sum_{\mathbf{k}} \frac{|g_{\mathbf{k}}|^2}{\omega_{\mathbf{k}}^2} [1 - \cos(\mathbf{k} \cdot \mathbf{d})] \coth(\beta\omega_{\mathbf{k}}/2) \right], \quad (6)$$

is the expectation value of the bath operators with respect to the free Hamiltonian: $B = \langle B_{\pm} \rangle_{H_0}$. The donor-acceptor separation is given by $\mathbf{d} = \mathbf{r}_1 - \mathbf{r}_2$.

To calculate the renormalization factor, we take the continuum limit to convert the summation in Eq. (6) into an integral. Defining the bath spectral density $J(\omega) = \sum_{\mathbf{k}} |g_{\mathbf{k}}|^2 \delta(\omega - \omega_{\mathbf{k}})$, which contains information regarding both the density of oscillators in the bath with a given frequency, and also how strongly those oscillators interact with the donor-acceptor pair, and assuming a linear, isotropic dispersion relation, we find

$$B = \exp \left[- \int_0^{\infty} \frac{J(\omega)}{\omega^2} [1 - F_D(\omega, d)] \coth(\beta\omega/2) d\omega \right]. \quad (7)$$

Here, $\beta = 1/k_B T$ is the inverse temperature, while the function $F_D(\omega, d)$ captures the degree of spatial correlation in the bath fluctuations seen at each site, and is dependent upon the dimensionality of the system-bath interaction ($D = 1, 2, 3$) (Refs. 33,53,67). In one dimension $F_1(\omega, d) = \cos(\omega d/c)$, with c the bosonic excitation speed, in two dimensions $F_2(\omega, d) = J_0(\omega d/c)$, where $J_0(x)$ is a Bessel function of the first kind, and in three dimensions $F_3(\omega, d) = \text{sinc}(\omega d/c)$. In all cases $F_D(\omega, d) \rightarrow 1$ as $d \rightarrow 0$, that is, when the donor and acceptor are at the same position, bath fluctuations are perfectly correlated, and the energy transfer strength is not renormalized ($V_R \rightarrow V$). In fact, in this limit dissipative processes are entirely suppressed (provided $|g_{\mathbf{k}}^{(1)}| = |g_{\mathbf{k}}^{(2)}|$) and energy transfer remains coherent for all times and in all parameter regimes in our model (the single-excitation subspace is then decoherence-free^{68,69}). In two and three dimensions, as $d \rightarrow \infty$, $F_D(\omega, d) \rightarrow 0$, and the renormalization takes on the value that would be obtained by considering separate, completely uncorrelated baths surrounding the donor and acceptor. In the following, we shall characterize the degree of correlation in terms of the dimensionless parameter $\mu = c/\omega_0 d$, where ω_0 is a typical bath frequency scale [see Eq. (29) below]. We therefore have $\mu = 0$ in the absence of correlations, $\mu < 1$ for weak correlations, and $\mu > 1$ for strong correlations.

B. Markovian master equation

Having identified a perturbation term by transforming our Hamiltonian to the polaron representation, we can now construct a master equation describing the evolution of the donor-acceptor pair reduced density operator ρ up to second order in H_I . We employ a standard Born-Markov approach, which yields a polaron frame, interaction picture master equation of the form²⁵

$$\frac{\partial \tilde{\rho}(t)}{\partial t} = - \int_0^{\infty} d\tau \text{tr}_B \{ [\tilde{H}_I(t), [\tilde{H}_I(t - \tau), \tilde{\rho}(t) \otimes \rho_B]] \}, \quad (8)$$

where tildes indicate operators in the interaction picture, $\tilde{O}(t) = e^{iH_0 t} O e^{-iH_0 t}$, and tr_B denotes a trace over the bath degrees of freedom. In deriving Eq. (8) we have assumed (i) factorizing initial conditions for the joint system-bath density operator within the polaron frame, $\chi(0) = \rho(0) \otimes \rho_B$, with

$\rho_B = e^{-\beta H_B} / \text{tr}_B(e^{-\beta H_B})$ being a thermal equilibrium state of the bath; (ii) that by construction the interaction is weak in the polaron frame so that we may factorize the joint density operator as $\tilde{\chi}(t) = \tilde{\rho}(t) \rho_B$ at all times; (iii) that the time scale on which the donor-acceptor system evolves appreciably in both the Schrödinger and interaction pictures is large compared to the bath memory time τ_B . Since, for the spectral density we shall consider below, $\tau_B \sim 1/\omega_c$, where ω_c is a high-frequency cutoff [see Eq. (29)], this is not too restrictive, as we must keep $V < \omega_c$ anyway in order for the polaron theory to work well. We note that interesting non-Markovian and nonequilibrium bath effects have been explored in the polaron formalism in Refs. 12 and 13.

Inserting Eq. (4) into Eq. (8), and moving back into the Schrödinger picture, we arrive at our Markovian master equation describing the energy transfer dynamics within the single-excitation subspace, and written in the polaron frame as

$$\begin{aligned} \frac{\partial \rho(t)}{\partial t} = & -i[(\epsilon/2)\sigma_z + V_R \sigma_x, \rho(t)] \\ & - V^2 \int_0^{\infty} d\tau \{ [\sigma_x, \tilde{\sigma}_x(-\tau) \rho(t)] \Lambda_{xx}(\tau) \\ & + [\sigma_y, \tilde{\sigma}_y(-\tau) \rho(t)] \Lambda_{yy}(\tau) + \text{H.c.} \}, \quad (9) \end{aligned}$$

where H.c. denotes Hermitian conjugation. The effect of the bath is now contained within the correlation functions $\Lambda_{ll}(\tau) = \langle \tilde{B}_l(\tau) \tilde{B}_l(0) \rangle_{H_0}$, which are given explicitly by

$$\Lambda_{xx}(\tau) = (B^2/2)(e^{\phi(\tau)} + e^{-\phi(\tau)} - 2), \quad (10)$$

$$\Lambda_{yy}(\tau) = (B^2/2)(e^{\phi(\tau)} - e^{-\phi(\tau)}), \quad (11)$$

where

$$\begin{aligned} \phi(\tau) = & 2 \int_0^{\infty} d\omega \left\{ \frac{J(\omega)}{\omega^2} [1 - F_D(\omega, d)] \right. \\ & \left. \times [\cos \omega \tau \coth(\beta\omega/2) - i \sin \omega \tau] \right\}. \quad (12) \end{aligned}$$

Notice that the phonon propagator $\phi(\tau)$ is correlation dependent due to the factor $[1 - F_D(\omega, d)]$, and so clearly the dissipative effect of the bath will be dependent upon the degree of correlation also. For example, as $d \rightarrow 0$, $F_D(\omega, d) \rightarrow 1$, and the dissipative contribution to Eq. (9) vanishes, as anticipated earlier.

C. Evolution of the Bloch vector

We solve our master equation in terms of the Bloch vector, defined as $\boldsymbol{\alpha} = (\alpha_x, \alpha_y, \alpha_z)^T = (\langle \sigma_x \rangle, \langle \sigma_y \rangle, \langle \sigma_z \rangle)^T$. As we are working exclusively in the single-excitation subspace, α_x and α_y describe the coherences between the states $|1\rangle \equiv |XG\rangle$ and $|2\rangle \equiv |GX\rangle$, while α_z captures the donor-acceptor population transfer dynamics generated by the coupling V .

Though Eq. (9) is written in the Schrödinger picture, it is still in the polaron frame, and so we must determine how expectation values in the polaron frame are related to those in the original, or ‘‘lab’’ frame. We can see this by writing $\alpha_i = \text{tr}_{S+B}(\sigma_i \chi_L(t)) = \text{tr}_{S+B}[\sigma_i e^{-S} \chi(t) e^S] = \text{tr}_{S+B}[e^S \sigma_i e^{-S} \rho(t) \rho_B]$, where $\chi_L(t) = e^{-S} \chi(t) e^S$ is the lab frame total density operator, and we have made use of the Born approximation in the polaron frame to write

$\chi(t) = \rho(t)\rho_B$. Since $e^S\sigma_x e^{-S} = |2\rangle\langle 1|B_- + |1\rangle\langle 2|B_+$, $e^S\sigma_y e^{-S} = i(|2\rangle\langle 1|B_- - |1\rangle\langle 2|B_+)$, and $e^S\sigma_z e^{-S} = \sigma_z$, this implies that the lab Bloch vector elements are $\alpha_i = B\alpha_{iP}$, for $i = x, y$, and $\alpha_z = \alpha_{zP}$, where α_{iP} is an expectation value in the polaron frame: $\alpha_{iP} = \text{Tr}_S[\sigma_i \rho(t)]$. Alternatively, we can define a matrix L which maps the polaron frame Bloch vector (α_P) to its lab frame counterpart (α): $\alpha = L \cdot \alpha_P$, where $L = \text{diag}(B, B, 1)$.

Working in terms of the Bloch vector, we arrive at an equation of motion of the form

$$\dot{\alpha}(t) = M \cdot \alpha(t) + \mathbf{b}. \quad (13)$$

In the following, we shall often be interested in determining whether the energy transfer dynamics is predominantly coherent or incoherent. It is then helpful to write Eq. (13) as

$$\dot{\alpha}'(t) = M \cdot \alpha'(t), \quad (14)$$

with $\alpha'(t) = \alpha(t) - \alpha(\infty)$, where $\alpha(\infty) = -M^{-1} \cdot \mathbf{b}$ is the steady state. This makes clear that the nature of the energy transfer process lies solely in the matrix M , while the inhomogeneous term \mathbf{b} is needed only in determining the steady state.

Equipped with the eigensystem of M , we may determine the corresponding time evolution as follows: an eigenvector of M , say \mathbf{m}_i , has equation of motion $\dot{\mathbf{m}}_i = q_i \mathbf{m}_i$, where q_i is the corresponding eigenvalue. Its subsequent evolution then has the simple exponential form $\mathbf{m}_i(t) = \mathbf{m}_i e^{q_i t}$. More generally, we can say that any initial state $\alpha'(0)$ will have the subsequent evolution

$$\alpha'(t) = \sum_{i=1}^3 a_i \mathbf{m}_i e^{q_i t}, \quad (15)$$

where the coefficients a_i are determined by the initial conditions [i.e., the solutions of $\alpha'(0) = \sum_i a_i \mathbf{m}_i$]. The solution to the full inhomogeneous equation is then found simply by the addition of the steady state $\alpha(t) = \alpha'(t) + \alpha(\infty)$.

III. RESONANT ENERGY TRANSFER

We start by considering the important special case of resonant energy transfer, in which the interplay of coherent and incoherent effects is particularly pronounced. As we shall see, in this situation it is relatively straightforward to derive a strict criterion governing whether or not we expect the energy transfer dynamics to be able to display signatures of coherence.^{11,51} Hence, resonant conditions provide a natural situation in which to begin to understand, for example, the role of bath spatial correlations^{9,11,33,53,56-60} or the range of the bath frequency distribution in determining the nature of the energy transfer process.

Setting the donor-acceptor energy mismatch to zero ($\epsilon = 0$), we find from Eq. (9) dynamics generated by an expression of the form $\dot{\alpha} = M_R \cdot \alpha + \mathbf{b}_R$, with

$$M_R = \begin{pmatrix} -(\Gamma_z - \Gamma_y) & 0 & 0 \\ 0 & -\Gamma_y & -2BV_R \\ 0 & B^{-1}(2V_R + \lambda_3) & -\Gamma_z \end{pmatrix}, \quad (16)$$

and $\mathbf{b}_R = (-B\kappa_x, 0, 0)^T$, where

$$\Gamma_y = 2V^2\gamma_{xx}(0), \quad (17)$$

$$\Gamma_z = V^2[\gamma_{yy}(2V_R) + \gamma_{yy}(-2V_R)] + 2V^2\gamma_{xx}(0), \quad (18)$$

$$\lambda_3 = 2V^2[S_{yy}(2V_R) - S_{yy}(-2V_R)], \quad (19)$$

$$\kappa_x = V^2[\gamma_{yy}(2V_R) - \gamma_{yy}(-2V_R)]. \quad (20)$$

The rates and energy shifts are related to the response functions

$$K_{ii}(\omega) = \int_0^\infty d\tau e^{i\omega\tau} \Lambda_{ii}(\tau) = \frac{1}{2}\gamma_{ii}(\omega) + iS_{ii}(\omega), \quad (21)$$

such that

$$\gamma_{ii}(\omega) = 2\text{Re}[K_{ii}(\omega)] = \int_{-\infty}^{+\infty} d\tau e^{i\omega\tau} \Lambda_{ii}(\tau), \quad (22)$$

and $S_{ii}(\omega) = \text{Im}[K_{ii}(\omega)]$.

The resonant steady state is straightforwardly found to be

$$\alpha_x(\infty) = -B \tanh(\beta V_R), \quad (23)$$

while $\alpha_y(\infty) = \alpha_z(\infty) = 0$. Notice that while this is similar in form to the steady state that would be obtained from a weak system-bath coupling treatment,^{25,70} $\alpha_x(\infty)$ is determined here by V_R , rather than the original coupling V , and there is also an extra factor of B suppressing its magnitude.

The eigenvalues of M_R are given by $q_1 = \Gamma_y - \Gamma_z$ and $q_2 = q_3^* = -(1/2)(\Gamma_y + \Gamma_z + i\xi_R)$. Thus, referring to Eq. (15), we see that

$$\xi_R = \sqrt{8V_R(2V_R + \lambda_3) - (\Gamma_z - \Gamma_y)^2}, \quad (24)$$

determines whether or not any coherence exists within the energy transfer dynamics. Considering the initial state $\alpha(0) = (0, 0, 1)^T$, corresponding to the excitation of the donor, $\rho(0) = |1\rangle\langle 1| = |XG\rangle\langle XG|$, we find analytical forms for the evolution of the Bloch vector components

$$\alpha_x(t) = -B \tanh(\beta V_R)(1 - e^{-(\Gamma_y - \Gamma_z)t}), \quad (25)$$

$$\alpha_y(t) = -\frac{2BV_R}{\xi_R} e^{-(\Gamma_y + \Gamma_z)t/2} \sin\left(\frac{\xi_R t}{2}\right), \quad (26)$$

$$\alpha_z(t) = e^{-(\Gamma_y + \Gamma_z)t/2} \left[\cos\left(\frac{\xi_R t}{2}\right) + \frac{\Gamma_y - \Gamma_z}{\xi_R} \sin\left(\frac{\xi_R t}{2}\right) \right]. \quad (27)$$

An inspection of Eqs. (24) and (27) allows us to identify a crossover from coherent to incoherent motion in the energy transfer dynamics as the point at which oscillations in the population difference vanish¹¹

$$(\Gamma_z - \Gamma_y)^2 = 8V_R(2V_R + \lambda_3). \quad (28)$$

For $(\Gamma_z - \Gamma_y)^2 < 8V_R(2V_R + \lambda_3)$, ξ_R is real and both the population difference and coherence α_y describe damped oscillations, while for $(\Gamma_z - \Gamma_y)^2 \geq 8V_R(2V_R + \lambda_3)$, ξ_R is either zero or imaginary, with the resulting dynamics then being entirely incoherent.

To further analyze the behavior of $\alpha_z(t)$, and the conditions for which the boundary defined by Eq. (28) is crossed, we now take a specific form for the system-bath spectral density. For a large enough bath we may approximate $J(\omega) = \sum_{\mathbf{k}} |g_{\mathbf{k}}|^2 \delta(\omega_{\mathbf{k}} - \omega)$ as a smooth function of ω . In this work we consider a spectral density of the form

$$J(\omega) = \alpha \frac{\omega^3}{\omega_0^2} e^{-\omega/\omega_c}, \quad (29)$$

where α is a dimensionless quantity capturing the strength of the system-bath interaction, and ω_0 is a typical frequency of bosons in the bath, which sets an overall energy scale. The cubic frequency dependence in Eq. (29) is typical, for example, in describing dephasing due to coupling to acoustic phonons,^{51,72} but can also be used to elucidate the behavior in which we are interested in general.⁵³ The cutoff frequency ω_c is needed to ensure that vacuum contributions remain finite, and is related to parameters specific to the particular physical system one wishes to model. The inverse cutoff frequency also sets a typical relaxation time scale for the bath.²⁵

To illustrate the dynamics and crossover behavior in the resonant case, in Fig. 1 we plot the population difference (α_z) as a function of the scaled time $\omega_0 t$ for a range of temperatures, showing the transition from coherent to incoherent transfer as the temperature is increased. In this plot, and all the following, we consider three-dimensional coupling, $F_3(\omega, d) = \text{sinc}(\omega d/c)$. The role of bath spatial correlations in protecting coherence can be seen in Fig. 2, where we again plot the evolution of the population difference (the insets show the corresponding coherence α_y), this time for representative intermediate and high-temperature cases. The different plots in Fig. 2 correspond to zero correlations, characterized by $\mu = c/\omega_0 d = 0$ ($d \rightarrow \infty$, top), weak correlations $\mu = 0.5$ (middle), and strong correlations $\mu = 2$ (bottom).⁷¹ Progressing from the uppermost plot to the lowest, we clearly see that an increase in correlation strength prolongs the time scale over which oscillations in both the population difference

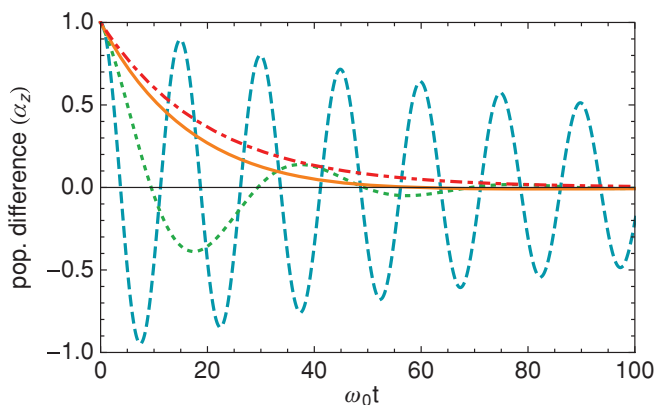


FIG. 1. (Color online) Population difference as a function of scaled time $\omega_0 t$ for temperatures of $k_B T/\omega_0 = 1$ (blue dashed curve), $k_B T/\omega_0 = 5$ (green dotted curve), $k_B T/\omega_0 = 12$ (orange solid curve), and $k_B T/\omega_0 = 20$ (red dot-dashed curve). Parameters: $\alpha = 0.05$, $V/\omega_0 = 0.5$, $\omega_c/\omega_0 = 4$, $\epsilon = 0$, and $\mu = c/\omega_0 d = 0.5$.

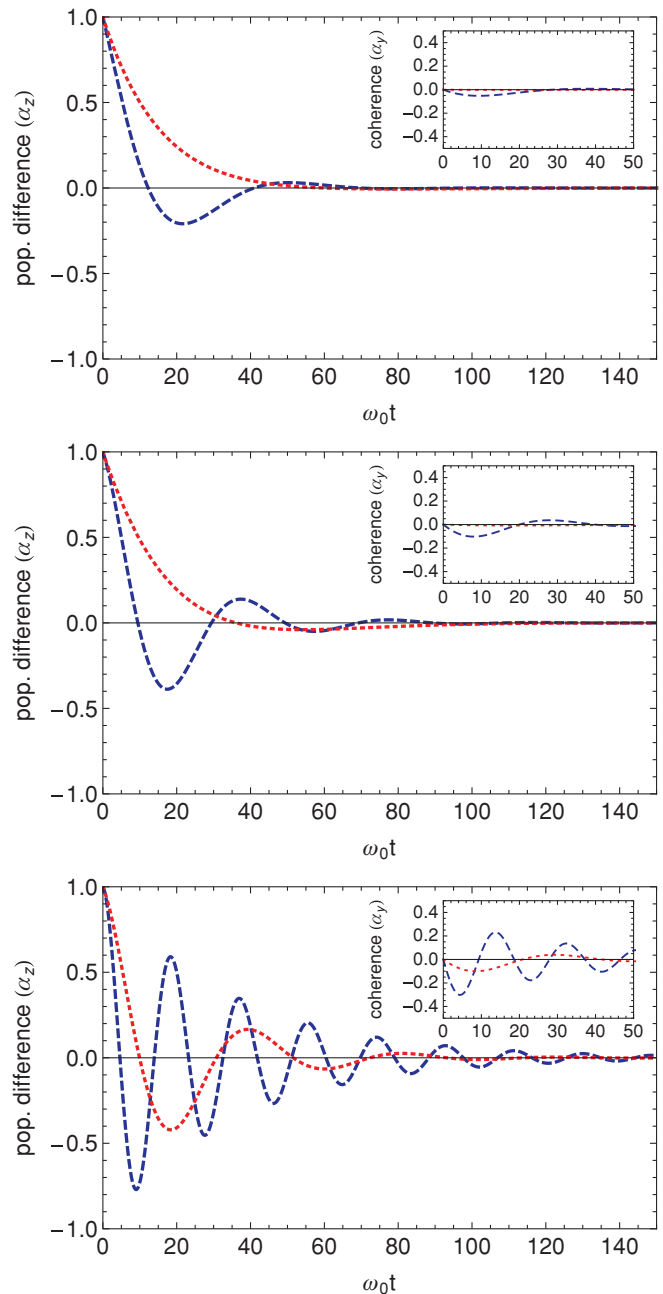


FIG. 2. (Color online) Population difference as a function of scaled time $\omega_0 t$ for temperatures of $k_B T/\omega_0 = 5$ (blue dashed curves) and $k_B T/\omega_0 = 10$ (red dotted curves), and for separations corresponding to no correlation, $\mu = c/\omega_0 d = 0$ (top), weak correlations, $\mu = 0.5$ (middle), and strong correlations $\mu = 2$ (bottom). The insets show the evolution of the corresponding coherence α_y . Parameters: $\alpha = 0.05$, $V/\omega_0 = 0.5$, and $\omega_c/\omega_0 = 4$.

and coherence persist. Moreover, by looking at the curves corresponding to the higher temperature (red, dotted), we can see that as the degree of correlation is increased from zero, the dynamics moves from a regime showing purely incoherent relaxation, to a regime which displays coherent oscillations *at the same temperature*. The increase in correlations is thus able to extend the region of parameter space which permits coherence,¹¹ as we shall now explore in greater detail.

A. Coherent to incoherent transition

We now return our attention to the crossover from coherent to incoherent transfer, defined by Eq. (28). Intuitively, we might expect the dynamics in the low-temperature (or weak-coupling) regime to be coherent; for example, in Fig. 1 incoherent relaxation only occurs in the high-temperature limit. If we therefore assume that the crossover itself occurs in the high-temperature regime, it is possible to derive an analytic expression governing the crossover temperature by approximating the rates Γ_y and Γ_z . Details of this approximation, and its range of validity, can be found in the Appendix. Generally, for high enough temperatures and/or strong enough system-bath coupling (such that $\beta V_R \ll 1$) we can approximate $\gamma_{xx}(\eta) \approx \gamma_{yy}(\eta) \approx \gamma_{yy}(0)$ in Γ_y and Γ_z , where

$$\gamma_{yy}(0) \approx \frac{\beta B^2 e^{\phi_0 C_0(x,y)}}{2\sqrt{\pi} C_2(x,y) \phi_0}, \quad (30)$$

with $\phi_0 = 2\pi^2 \alpha / \omega_0^2 \beta^2$, $x = \pi d / c \beta$ and $y = \omega_c \beta$. The functions $C_0(x,y)$ and $C_2(x,y)$ are given by Eqs. (A7) and (A8), and the renormalization factor B by the product of Eqs. (A14) and (A15). If we further assume that the energy shift λ_3 vanishes in the high-temperature limit, Eq. (28) reduces to

$$(\Gamma_z - \Gamma_y) = 4V_R, \quad (31)$$

and we arrive at the expression

$$\left(\frac{k_B T}{\omega_0}\right)^2 = \frac{V}{\omega_0} \frac{B e^{\phi_0 C_0(x,y)}}{4\sqrt{2\pi^3 \alpha} C_2(x,y)}, \quad (32)$$

with solution, T_c , giving the crossover temperature separating the coherent and incoherent regimes.

The dependence of T_c on the various parameters involved in the problem is not straightforward, owing to the temperature dependence in the renormalization factor B , in the functions C_0 and C_2 , and in ϕ_0 . In fact, there are three distinct and important temperature scales which determine when coherent or incoherent processes dominate: $T_0 = \omega_0 / (\sqrt{2\pi} \alpha k_B)$, which depends upon the system-bath coupling strength; $T_x = c / d \pi k_B$, which arises due to the fluctuation correlations and becomes unimportant in the uncorrelated case ($T_x \rightarrow 0$ as $d \rightarrow \infty$); and $T_y = \omega_c / k_B$, dependent upon the cutoff frequency, and irrelevant in the scaling limit ($y \rightarrow \infty$). Hence, changes in any of α , d , or ω_c can have an effect on the crossover temperature. For example, the main part of Fig. 3 shows the solution to Eq. (32) (i.e., the crossover temperature T_c) as a function of the dimensionless cutoff frequency ω_c / ω_0 . A calculation using Eq. (28) with the full rates, and including λ_3 , is also shown for comparison. The three pairs of curves correspond to increasing levels of correlation, ordered as indicated. We see that, except for small ω_c / ω_0 in the case $\mu = 0$ where λ_3 becomes important, solutions to Eq. (32) give an excellent approximation to the crossover temperature calculated using the full rates. This confirms that the coherent-incoherent crossover does indeed occur in the high-temperature (multiphonon) regime, and consequently could not be captured by a weak system-bath coupling treatment.

As the cutoff frequency is increased from its minimum value, the crossover temperature begins to decrease. This behavior can be understood qualitatively by examining

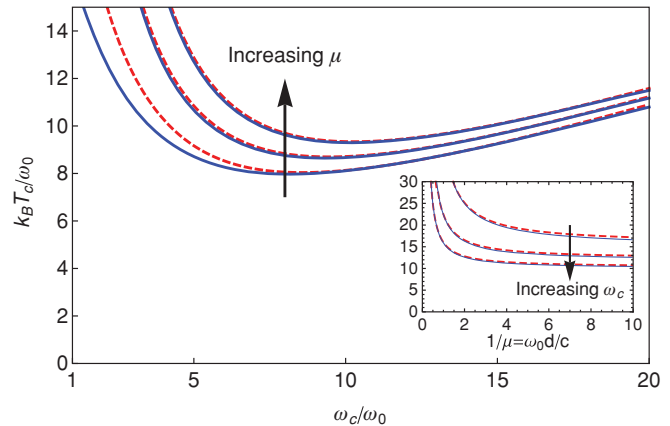


FIG. 3. (Color online) Crossover temperature separating the coherent and incoherent regimes against cutoff frequency, for levels of correlation given by $\mu = c/\omega_0 d = 0$, $\mu = 0.5$ and $\mu = 1$, increasing as shown. The solid blue curves have been calculated from Eq. (28) (using the full rates), while the dashed red curves are solutions to the high-temperature approximation, Eq. (32). The inset shows the dependence on the level of correlation ($1/\mu = \omega_0 d/c$) for different cutoffs, $\omega_c/\omega_0 = 2$, $\omega_c/\omega_0 = 3$, and $\omega_c/\omega_0 = 4$, again increasing as shown. Parameters: $\alpha = 0.05$ and $V/\omega_0 = 0.5$.

Eq. (31), and considering the competition this condition captures between the rate $\Gamma_z - \Gamma_y$ and the coherent interaction V_R in defining the nature of the dynamics. Larger values of the cutoff frequency correspond to smaller values of the renormalized interaction strength V_R [see, e.g., Eq. (A14)], while the rates Γ_y and Γ_z vary less strongly with ω_c in this regime. Thus, increasing ω_c from its minimum value decreases V_R , and therefore reduces the range of temperatures for which $4V_R > \Gamma_z - \Gamma_y$ and coherent transfer can take place. Thus, the crossover temperature falls. Physically, this can be understood by noting that as the cutoff frequency is increased, so too is the effective frequency range and peak magnitude of the system-bath interaction, characterized by the spectral density [Eq. (29)]. Hence, increasing from small ω_c / ω_0 , the environment begins to exert an enhanced influence on the system behavior, and so coherent dynamics no longer survives to such high temperatures. As ω_c continues to increase, however, we see the crossover temperature then begins to rise. The renormalization factor B tends to zero with increasing ω_c and here becomes the dominating quantity, thus causing the rate $\Gamma_z - \Gamma_y \sim \mathcal{O}(B^2)$ to vanish faster than the renormalized donor-acceptor coupling $V_R = BV$.

The interplay between the size of ω_c and the level of spatial correlation is best understood by considering the inset of Fig. 3. For all curves shown the crossover temperature increases as the distance d is reduced since the level of correlation μ increases correspondingly. As we have seen previously in Fig. 2, stronger correlations allow coherent dynamics to be observed at higher temperatures; since environmental effects are suppressed, so the crossover temperature T_c must rise. This behavior can be attributed to an increase in the renormalized interaction strength, V_R , in relation to the rate $\Gamma_z - \Gamma_y$, this time with variations in the correlation level μ . Interestingly, as the cutoff frequency is increased up to $\omega_c / \omega_0 = 4$ (lowest curve), we see that not only does the crossover temperature decrease, but also

that the degree of correlation necessary to show a marked rise in T_c increases. As can be seen by comparing the separation between the different curves in the main part of the figure, increasing the cutoff frequency tends to suppress the extent to which correlations are able to protect coherence in the system. This tallies with the dynamics shown in Fig. 2, for which $\omega_c/\omega_0 = 4$, and correlations as high as $\mu = 2$ were needed before a significant change in behavior was seen. Finally, since the renormalization factor B tends to a constant nonzero value as the correlations vanish at large d (as opposed to $B \rightarrow 0$ as $\omega_c \rightarrow \infty$), the dependence of the crossover temperature on μ is monotonic, in contrast to its dependence on ω_c .

IV. OFF RESONANCE

It is often the case, in practice, that the donor and acceptor will have different excited state energies, $\epsilon_1 - \epsilon_2 = \epsilon \neq 0$, and so we now turn our attention to energy transfer dynamics under off-resonant conditions. Regarding the coherent to incoherent transition, in the resonant case we were able to identify this point with a pair of conjugate eigenvalues converging on the real axis, thus changing oscillatory terms into relaxation. We might hope that in the off-resonant case we are able to establish a similar crossover criterion, and again use this to investigate the effects of bath correlations and the cutoff frequency. However, we shall see that such an identification is less straightforward in the off-resonant regime.

We first present the full Bloch equations describing the evolution of our donor-acceptor pair for arbitrary energy mismatch. As in the resonant case, we have an equation of motion of the form $\dot{\alpha} = M \cdot \alpha + \mathbf{b}$, but now the matrix M is given by

$$M = \begin{pmatrix} -\Gamma_x & -(\epsilon + \lambda_1) & 0 \\ (\epsilon + \lambda_2) & -\Gamma_y & -2BV_R \\ B^{-1}\zeta & B^{-1}(2V_R + \lambda_3) & -\Gamma_z \end{pmatrix}, \quad (33)$$

with $\mathbf{b} = (-B\kappa_x, -B\kappa_y, -\kappa_z)^T$. The rates become

$$\Gamma_x = V^2[\gamma_{yy}(\eta) + \gamma_{yy}(-\eta)], \quad (34)$$

$$\Gamma_y = 2V^2 \left(\frac{4V_R^2}{\eta^2} \gamma_{xx}(0) + \frac{\epsilon^2}{2\eta^2} [\gamma_{xx}(\eta) + \gamma_{xx}(-\eta)] \right), \quad (35)$$

with $\Gamma_z = \Gamma_x + \Gamma_y$, and the energy shifts

$$\lambda_1 = \frac{2V^2\epsilon}{\eta} [S_{yy}(\eta) - S_{yy}(-\eta)], \quad (36)$$

$$\lambda_2 = \frac{2V^2\epsilon}{\eta} [S_{xx}(\eta) - S_{xx}(-\eta)], \quad (37)$$

$$\lambda_3 = \frac{4V^2V_R}{\eta} [S_{yy}(\eta) - S_{yy}(-\eta)]. \quad (38)$$

The remaining quantities are

$$\zeta = \frac{4V^2V_R\epsilon}{\eta^2} \left(\gamma_{xx}(0) - \frac{1}{2} [\gamma_{xx}(\eta) + \gamma_{xx}(-\eta)] \right), \quad (39)$$

$$\kappa_x = \frac{2V^2V_R}{\eta} [\gamma_{yy}(\eta) - \gamma_{yy}(-\eta)], \quad (40)$$

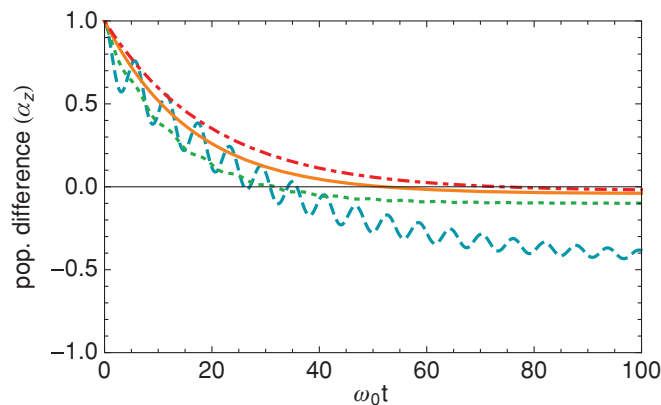


FIG. 4. (Color online) Population difference for an off-resonant donor-acceptor pair as a function of scaled time $\omega_0 t$. Temperatures $k_B T/\omega_0 = 1$ (blue dashed curve), $k_B T/\omega_0 = 5$ (green dotted curve), $k_B T/\omega_0 = 12$ (orange solid curve), and $k_B T/\omega_0 = 20$ (red dot-dashed curve) are shown. Parameters: $\alpha = 0.05$, $V/\omega_0 = 0.5$, $\epsilon/\omega_0 = 1$, $\omega_c/\omega_0 = 4$, and $\mu = 0.5$.

$$\kappa_y = \frac{8V^2V_R\epsilon}{\eta^2} \left(S_{xx}(0) - \frac{1}{2} [S_{xx}(\eta) + S_{xx}(-\eta)] \right), \quad (41)$$

$$\kappa_z = \frac{V^2\epsilon}{\eta} \{ [\gamma_{xx}(\eta) - \gamma_{xx}(-\eta)] + [\gamma_{yy}(\eta) - \gamma_{yy}(-\eta)] \}. \quad (42)$$

Here, $\eta = \sqrt{\epsilon^2 + 4V_R^2}$ is the system Hamiltonian eigenstate splitting in the polaron frame.

To exemplify the dynamics generated by the full off-resonant Bloch equations, in Fig. 4 we plot the evolution of the population difference in the case of donor-acceptor energy mismatch, $\epsilon = 2V$. By comparison of Fig. 1 (plotted in the resonant case) and Fig. 4, we see that the presence of a substantial energy mismatch causes the low-temperature population oscillations to increase in frequency but decrease markedly in amplitude, such that for $k_B T/\omega_0 = 5$ oscillations are now almost imperceptible. We also see that the population difference tends to a nonzero steady state at low temperatures, as we might expect from simple thermodynamic arguments, since the states $\alpha_z = 1$ and $\alpha_z = -1$ now have different energies. As the temperature is raised, however, the dynamics still looks to be approaching that shown in the resonant case of Fig. 1.

A. Correlated fluctuations

Let us now look at the effect of correlated fluctuations in the off-resonant dynamics. Though analysis of the full Bloch equations is now more complicated than in the resonant limit, we should still expect changes in the level of donor-acceptor fluctuation correlation to have a qualitatively similar effect on the transfer process as outlined in Sec. III. To illustrate that this is indeed the case, in Fig. 5 we plot the donor-acceptor population dynamics under off-resonant conditions for three different levels of fluctuation correlation (increasing from top to bottom). Just as we found in the resonant case of Fig. 2, an increase in correlations enhances the lifetime of coherence

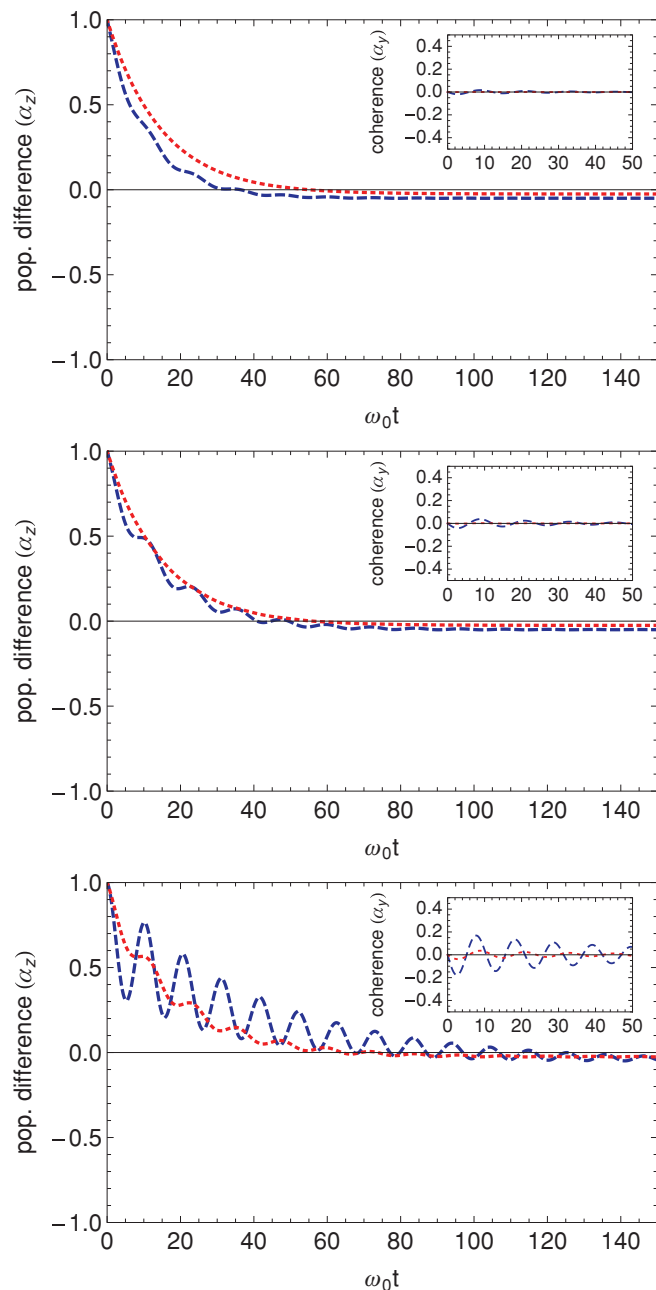


FIG. 5. (Color online) Population difference as a function of scaled time $\omega_0 t$ for temperatures of $k_B T / \omega_0 = 5$ (blue dashed curve) and $k_B T / \omega_0 = 10$ (red dotted curve), and for separations corresponding to no fluctuation correlations, $\mu = c / \omega_0 d = 0$ (top), weak correlations, $\mu = 0.5$ (middle), and strong correlations $\mu = 2$ (bottom). The insets show the evolution of the corresponding coherence α_y . Parameters: $\alpha = 0.05$, $V / \omega_0 = 0.5$, $\epsilon / \omega_0 = 0.5$, and $\omega_c / \omega_0 = 4$.

present in the energy transfer process, and can even move the dynamics from a high temperature (or strong-coupling) predominantly incoherent regime to an *effective* low temperature (or weak-coupling) regime displaying pronounced coherent oscillations. In addition, in the off-resonant case stronger correlations also serve to amplify the coherent contribution to the full energy transfer dynamics (made up of distinct coherent and incoherent parts, as we shall show below) since

the renormalized interaction strength V_R increases in relation to the energy mismatch ϵ .

In an effort to put these qualitative observations on a more quantitative footing we could analyze the eigensystem of the full off-resonant M [Eq. (33)], in a similar manner to the resonant case. However, finding the eigensystem of M is now far less straightforward and analytical solutions to the full Bloch equations are consequently lengthy, and therefore of little direct use in gaining an understanding of the behavior seen in Figs. 4 and 5. The rest of this section is thus devoted to deriving simplified expressions for the energy transfer dynamics in two important limits: (i) weak-coupling, or strong correlations, where coherent dynamics can dominate, and (ii) high temperatures (and weak correlations), where the dynamics is similar in both the resonant and off-resonant cases. These expressions not only provide insight into the off-resonant behavior of the system and the effect of correlated fluctuations, but also serve to highlight the difficulty in now defining a simple crossover criterion, as was possible in the resonant case.

B. Weak-coupling (or strong correlation) limit

We begin by considering the weak system-bath coupling limit, which we obtain by expanding all relevant quantities to first order in $J(\omega)$. In fact, for strong enough fluctuation correlations this limit is attainable even if the system-bath coupling is not weak and/or the temperature not low due to the factor $[1 - F_D(\omega, d)]$ appearing in Eqs. (7) and (12). With reference to our expressions for the correlation functions [Eqs. (10) and (11)], we see that within this approximation $\Lambda_{xx}(\tau) \rightarrow 0$ while $\Lambda_{yy}(\tau)$ remains finite. We may then set to zero all rates and energy shifts which are functions of $\Lambda_{xx}(\tau)$ only in Eq. (33). This results in the far simpler form

$$M_W = \begin{pmatrix} -\Gamma_W & -(\epsilon + \lambda_1) & 0 \\ \epsilon & 0 & -2BV_R \\ 0 & B^{-1}(2V_R + \lambda_3) & -\Gamma_W \end{pmatrix}, \quad (43)$$

where the weak-coupling rate is given by⁷³

$$\Gamma_W = 4\pi \left(\frac{V_R}{\eta} \right)^2 J(\eta) [1 - F_D(\eta, d)] \coth(\beta\eta/2), \quad (44)$$

and the two energy shifts may be written $\lambda_1 = (\epsilon/\eta)\Lambda$ and $\lambda_3 = (2V_R/\eta)\Lambda$, with $\Lambda = 2V^2[S_{yy}(\eta) - S_{yy}(-\eta)]$. The inhomogeneous term becomes $\mathbf{b}_W = \{-B\kappa_x, 0, -(\epsilon/2V_R)\kappa_x\}^T$ in the same limit, which leads to the weak-coupling steady state values of

$$\alpha_x(\infty) = -\frac{2BV_R}{\eta} \tanh(\beta\eta/2), \quad (45)$$

$$\alpha_z(\infty) = -\frac{\epsilon}{\eta} \tanh(\beta\eta/2), \quad (46)$$

and $\alpha_y(\infty) = 0$. As in the resonant case (in which there was no weak-coupling approximation), this steady state has the same form as that expected from a standard weak-coupling approach, though with the replacement $V \rightarrow V_R$, and the extra factor of B suppressing the coherence $\alpha_x(\infty)$. As the energy mismatch increases in relation to V , the weak-coupling steady state therefore becomes increasingly localized in the

lower energy state $|2\rangle \equiv |GX\rangle$. Interestingly, this contrasts with the qualitatively incorrect form (at low temperatures at least) given by the noninteracting blip approximation (NIBA), $\alpha_z^{\text{NIBA}}(\infty) = -\tanh(\beta\epsilon/2)$ (Refs. 70,74), which predicts complete localization in the lower energy state at zero temperature, regardless of the size of ϵ/V . We should thus expect the present theory to fair far better than the NIBA for low temperatures (or weak coupling) in the off-resonant case, $\epsilon \neq 0$. The rate Γ_W given in Eq. (44) is also of the form expected from a weak-coupling treatment, though once more with the renormalization $V \rightarrow V_R$. In fact, such a replacement is sometimes made by hand in weak-coupling theories to provide agreement with numerics over a larger range of parameters,⁷³ though it arises naturally in the polaron formalism here. We can therefore conclude that, in addition to allowing for the exploration of multiphonon effects,^{11,12,51,55} the polaron master equation provides a rigorous way to explore the (single-phonon) weak-coupling regime for spectral densities of the type in Eq. (29).⁵¹

As before, to find the time evolution of α we evaluate the eigensystem of M_W and use Eq. (15). For the initial state $\alpha(0) = \{0, 0, 1\}^T$ we find population dynamics

$$\alpha_z(t) = \frac{\epsilon}{\eta} \left(\frac{\epsilon}{\eta} e^{-\Gamma_W t} - (1 - e^{-\Gamma_W t}) \tanh(\beta\eta/2) \right) + \frac{4V_R^2}{\eta^2} e^{-\frac{\Gamma_W t}{2}} \left[\cos\left(\frac{\xi_W t}{2}\right) - \frac{\Gamma_W}{\xi_W} \sin\left(\frac{\xi_W t}{2}\right) \right], \quad (47)$$

where the weak coupling oscillation frequency is given by

$$\xi_W = \sqrt{4\eta(\eta + \Lambda) - \Gamma_W^2}, \quad (48)$$

which we expect to be real to be consistent with our original expansion. The first term in Eq. (47), proportional to (ϵ/η) and present nowhere in the resonant case, describes incoherent relaxation toward the steady state value given by Eq. (46). The second term, proportional to $(V_R/\eta)^2$ and having a similar form to the resonant dynamics, describes damped oscillations with frequency ξ_W . Importantly, these oscillations have a temporal maximum amplitude of $4V_R^2/\eta^2 \leq 1$, compared to 1 in the resonant case. The effect of the energy mismatch in this limit is thus to suppress the amplitude of any oscillations in the population difference, while increasing their frequency due to the dependence of ξ_W on η in Eq. (48), exactly as seen in Fig. 4.

C. High temperature (or far from resonance) limit

At high temperatures and weak correlations, we find that the population dynamics appears to be relatively insensitive to the size of the energy mismatch. To investigate this effect in more detail, we shall now make a high-temperature (or strong system-bath coupling) approximation to the full energy transfer dynamics.

Specifically, we consider the regime $V_R/\epsilon \ll 1$. This limit can, in fact, be achieved in two possible ways. First, recalling that $V_R = BV$, we see that V_R can be made small by increasing the system-bath coupling strength or temperature, such that $B \ll 1$. Alternatively, if the donor-acceptor pair are far from resonance, the ratio V/ϵ will be small, and hence V_R/ϵ smaller still. Observing that the correlation functions given by Eqs. (10) and (11) are both proportional to B^2 , we can see that

all dissipative terms in the equation of motion, $\dot{\alpha} = M \cdot \alpha + \mathbf{b}$, are at least of order V_R^2 . We proceed by keeping only terms up to order $(V_R/\epsilon)^2$ in the full off-resonant M and \mathbf{b} . This allows us to set $\lambda_3, \zeta, \kappa_x$ and κ_y to zero, while the remaining quantities reduce to

$$\Gamma_y = V^2[\gamma_{xx}(\eta) + \gamma_{xx}(-\eta)], \quad (49)$$

$$\Gamma_z = V^2[\gamma_{xx}(\eta) + \gamma_{xx}(-\eta) + \gamma_{yy}(\eta) + \gamma_{yy}(-\eta)], \quad (50)$$

$$\lambda_1 = 2V^2[S_{yy}(\eta) - S_{yy}(-\eta)], \quad (51)$$

$$\lambda_2 = 2V^2[S_{xx}(\eta) - S_{xx}(-\eta)], \quad (52)$$

$$\kappa_z = V^2[\gamma_{xx}(\eta) - \gamma_{xx}(-\eta) + \gamma_{yy}(\eta) - \gamma_{yy}(-\eta)]. \quad (53)$$

Hence, in the high-temperature limit, Eq. (33) takes on the simpler form

$$M_{\text{HT}} = \begin{pmatrix} -(\Gamma_z - \Gamma_y) & -(\epsilon + \lambda_1) & 0 \\ (\epsilon + \lambda_2) & -\Gamma_y & -2BV_R \\ 0 & 2B^{-1}V_R & -\Gamma_z \end{pmatrix}, \quad (54)$$

while the inhomogeneous term reduces to $\mathbf{b}_{\text{HT}} = \{0, 0, -\kappa_z\}^T$. We then find the approximate steady-state population difference

$$\alpha_z(\infty) = - \left[1 + \frac{4V_R^2}{\epsilon^2} \left(\frac{\Gamma_y}{\Gamma_z} - 1 \right) \right] \tanh(\beta\eta/2), \quad (55)$$

valid up to second order in V_R/ϵ . For $V_R \ll \epsilon$, this steady-state is strongly localized in the low energy state $[\alpha_z(\infty) \approx -1]$ if $\epsilon \gg k_B T$, though for $\epsilon \ll k_B T$ thermal effects dominate and $\alpha_z(\infty) \approx 0$ as in the resonant case. Again, this behavior tallies with Fig. 4.

To obtain the corresponding population dynamics, we note in reference to Eq. (15) that the coefficients a_i , the eigenvectors \mathbf{m}_i , and the eigenvalues q_i will contain powers of our expansion parameter V_R/ϵ . Expanding both q_i and the products $a_i \mathbf{m}_i$ to second order, we find

$$\alpha_z(t) = e^{-\Gamma_z t} \left(1 - \frac{4V_R^2}{\epsilon^2} \right) + \frac{4V_R^2}{\epsilon^2} e^{-\Gamma_z t/2} \cos(\bar{\epsilon} t) - (1 - e^{-\Gamma_z t}) \tanh\left(\frac{\beta\eta}{2}\right) \left[1 + \frac{4V_R^2}{\epsilon^2} \left(\frac{\Gamma_y}{\Gamma_z} - 1 \right) \right], \quad (56)$$

where the shifted oscillation frequency is

$$\bar{\epsilon} = \epsilon + (1/2)(\lambda_1 + \lambda_2) + 2\epsilon(V_R/\epsilon)^2. \quad (57)$$

As in the weak-coupling case [Eq. (47)] the evolution of the donor-acceptor population difference consists of two contributions; incoherent relaxation toward the steady-state, and an oscillatory component with vanishing amplitude as $V_R/\epsilon \rightarrow 0$. The energy mismatch again serves to suppress oscillations in the population difference.

The most striking feature, however, of Eq. (56) is that there is an oscillatory component at frequency $\bar{\epsilon}$ at all. In the high-temperature limit, we might expect that this frequency would reach a point where it becomes imaginary and $\alpha_z(t)$ displays purely incoherent relaxation, as in the equivalent

resonant case. However, we can see that this is not the case since $\bar{\epsilon}$ is always real by definition. Furthermore, at very high temperatures $\bar{\epsilon} \rightarrow \epsilon$, and it therefore also remains finite. Eq. (56) thus highlights an important difference between the energy transfer dynamics in resonant and off-resonant situations. In the resonant case, as temperature is increased, the energy transfer process becomes less coherent through a reduction in oscillation frequency (i.e., V_R becomes small in comparison to $\Gamma_z - \Gamma_y$), eventually reaching a point at which population relaxes incoherently toward the steady state. In the off-resonant case, the transfer process becomes less coherent predominately through a reduction in oscillation amplitude. For high temperatures, an oscillatory component is still (in theory) present in the system, although it becomes ever more dominated by incoherent relaxation toward the steady-state population distribution, which depends upon the ratio $\epsilon/k_B T$. These features are clearly seen in Fig. 4.

Only to first order in V_R/ϵ do our expressions predict purely incoherent off-resonant population transfer:

$$\alpha_z(t) = e^{-\Gamma_z t} - (1 - e^{-\Gamma_z t}) \tanh(\beta\eta/2). \quad (58)$$

Let us also consider the evolution of α_y in the same limit:

$$\alpha_y(t) = -\frac{2BV_R}{\epsilon} e^{-(1/2)\Gamma_z t} \sin(\bar{\epsilon}t). \quad (59)$$

Hence, although the donor-acceptor population itself evolves entirely incoherently in this limit, the coherences may still perform oscillations due to the energy mismatch. To illustrate the difference in the transition to incoherent population transfer on- and off-resonance, in the main part of Fig. 6 we plot the evolution of the coherence $\alpha_y(t)$ in both cases. The parameters have been chosen such that the resonant dynamics is in the incoherent regime ($T > T_c$), hence the resonant α_y displays no oscillations [see Eq. (26)]. In accordance with Eq. (59), however, the introduction of an energy mismatch induces oscillations in the donor-acceptor coherence. While these oscillations have an almost negligible amplitude, this behavior

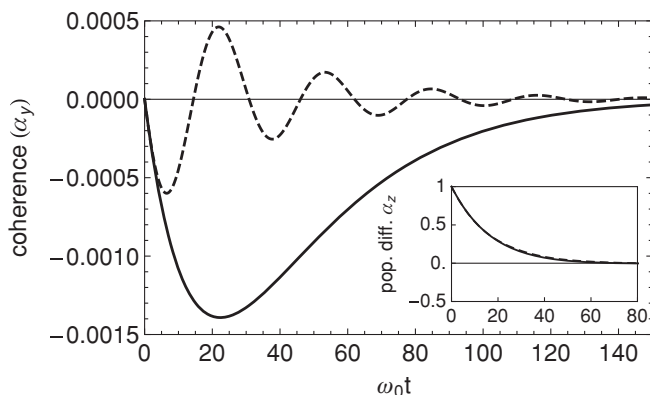


FIG. 6. Coherence (α_y) as a function of scaled time $\omega_0 t$ for resonant ($\epsilon = 0$, solid curve) and off-resonant ($\epsilon/\omega_0 = 0.2$, dashed curve) cases. The temperature, $k_B T/\omega_0 = 13$, is chosen to be above the relevant crossover T_c in the resonant case, such that the resonant dynamics is guaranteed to be incoherent. Parameters: $\alpha = 0.05$, $V/\omega_0 = 0.5$, $\omega_c/\omega_0 = 4$, and $\mu = 0.5$. The inset shows the corresponding population dynamics.

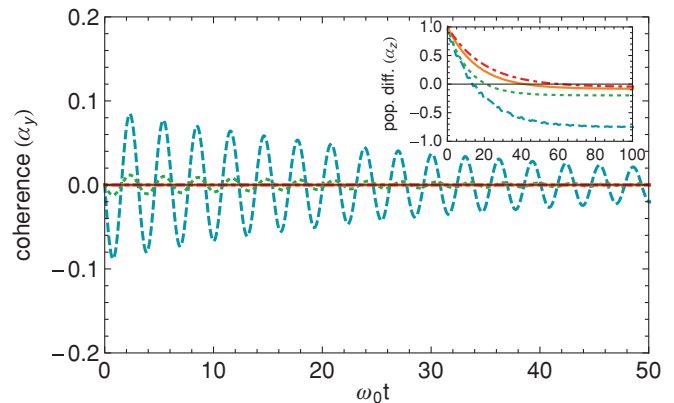


FIG. 7. (Color online) Coherence (α_y) as a function of scaled time $\omega_0 t$ for temperatures of $k_B T/\omega_0 = 1$ (blue dashed curve), $k_B T/\omega_0 = 5$ (green dotted curve), $k_B T/\omega_0 = 12$ (orange solid curve), and $k_B T/\omega_0 = 20$ (red dot-dashed curve). Parameters: $\alpha = 0.05$, $V/\omega_0 = 0.5$, $\omega_c/\omega_0 = 4$, $\epsilon/\omega_0 = 2$, and $\mu = 0.5$. The inset shows the corresponding population dynamics.

serves to illustrate the subtlety in defining a strict crossover from coherent to incoherent dynamics in the off-resonant case. In particular, despite the different forms of coherence behavior, the corresponding (essentially incoherent) population dynamics shown in the inset is almost indistinguishable in the two cases, even though there should still be a strongly suppressed coherent contribution in the off-resonant curve.

An alternative way to obtain oscillations of the coherence α_y in a regime of predominantly incoherent population transfer is to introduce a large energy mismatch (i.e., make V/ϵ small) at low temperature, as shown in Fig. 7. Here, for the lowest temperature considered the population relaxes toward its steady-state value with little sign of oscillation, while the coherence performs oscillations with a significant amplitude and considerable lifetime. This behavior is strongly suppressed, however, as temperature increases, such that $k_B T > \epsilon$.

V. SUMMARY

Motivated by recent experiments which suggest that quantum coherence can survive in energy transfer processes even under potentially adverse environmental conditions,¹⁻⁹ we have investigated various factors that determine the nature of the energy transfer dynamics in a model donor-acceptor pair. To do so, we used a polaron transform, Markovian master equation technique.¹¹ This formalism is attractive as it allows for exploration of both the low-temperature (or weak-coupling) and high-temperature (or strong-coupling) regimes, as well as reliable interpolation between these two limits, provided the ratio V/ω_c does not become too large.^{13,51,55} We are also able to consistently describe off-resonant effects, unlike in the NIBA,^{70,74} and the influence of bath correlations.

In the resonant case we identified a crossover temperature separating coherent and incoherent energy transfer. We found a nontrivial dependence of this temperature on both the degree of spatial correlation within the bath-induced fluctuations,

and also on the cutoff frequency of the bath spectral density. Smaller cutoff frequencies were found to enhance the extent to which bath spatial correlations are able to protect coherence in the system. The crossover generally occurs in a high-temperature limit where multiphonon effects dominate, and so could not be captured by a standard perturbative treatment of the system-bath interaction.

In the off-resonant case we found that coherent and incoherent regimes are less easily defined. In particular, for a sufficiently large energy mismatch between the donor and acceptor, coherence can, in theory, be present at all but infinite temperatures, albeit with an ever decreasing amplitude. However, using analytic expressions derived in various limits, we were able to characterize the off-resonant energy transfer process over much of the parameter space, illustrating the suppression of coherence in the population dynamics with increasing temperature or energy mismatch. We also showed that strong correlations have a qualitatively similar effect to the resonant case, protecting coherence in the transfer process.

While we have concentrated in this work on elucidating general features of donor-acceptor energy transfer dynamics using a simple model system, the insight we have gained could be relevant to a variety of systems. In addition to those already mentioned,¹⁻⁹ closely spaced pairs of semiconductor quantum dots could provide a solid-state implementation of the model studied here.⁷⁵ In particular, our polaron master equation theory provides a bridge between the weak⁷³ and strong⁷⁶ system-bath coupling approximations already explored in this context. It would also be interesting to analyze the energy transfer dynamics of larger donor-acceptor complexes within the polaron formalism⁷⁷ to see if further understanding of the interplay between coherent and incoherent processes in such systems could be obtained. Finally, it would be desirable to perform a thorough investigation of the regime of validity of the polaron approach by comparison to numerically exact techniques.⁷⁸

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APPENDIX : HIGH-TEMPERATURE RATES

Here we show how to obtain analytic approximations for the decoherence rates at high temperatures by use of a saddle point integration. Within our formalism there are two rates which need to be evaluated:

$$\gamma_{xx}(\eta) = \frac{B^2}{2} \int_{-\infty}^{+\infty} d\tau e^{i\tau\eta} (e^{\phi(\tau)} + e^{-\phi(\tau)} - 2), \quad (\text{A1})$$

$$\gamma_{yy}(\eta) = \frac{B^2}{2} \int_{-\infty}^{+\infty} d\tau e^{i\tau\eta} (e^{\tilde{\phi}(\tau)} - e^{-\tilde{\phi}(\tau)}), \quad (\text{A2})$$

where $\phi(\tau)$ is given by Eq. (12). With the appropriate manipulations⁵¹ it is possible to show that

$$\gamma_{xx}(\eta) = \frac{B^2}{2} e^{\beta\eta/2} \int_{-\infty}^{+\infty} d\tau e^{i\tau\eta} (e^{\tilde{\phi}(\tau)} + e^{-\tilde{\phi}(\tau)} - 2), \quad (\text{A3})$$

$$\gamma_{yy}(\eta) = \frac{B^2}{2} e^{\beta\eta/2} \int_{-\infty}^{+\infty} d\tau e^{i\tau\eta} (e^{\tilde{\phi}(\tau)} - e^{-\tilde{\phi}(\tau)}), \quad (\text{A4})$$

where now $\tilde{\phi}(\tau) = \tilde{\phi}(-\tau) = \phi(\tau - i\beta/2)$, and is given explicitly in integral form by

$$\tilde{\phi}(\tau) = 2 \int_0^{\infty} d\omega \frac{J(\omega)}{\omega^2} [1 - F_D(\omega, d)] \frac{\cos(\omega\tau)}{\sinh(\beta\omega/2)}. \quad (\text{A5})$$

Using a super-Ohmic form of spectral density, $J(\omega) = \alpha\omega^3\omega_0^{-2}e^{-\omega/\omega_c}$, and assuming system-bath coupling in three dimensions such that $F_D(\omega, d) = \text{sinc}(\omega d/c)$, allows $\tilde{\phi}(\tau)$ to be found analytically. We find $\tilde{\phi}(\tau) = \phi_0 C(x, y, \tau')$, where

$$\begin{aligned} C(x, y, \tau') = & \frac{-i}{2\pi x} \left[\psi \left(\frac{1}{2} + \frac{1}{y} - \frac{i}{\pi}(\tau' + x) \right) \right. \\ & - \psi \left(\frac{1}{2} + \frac{1}{y} - \frac{i}{\pi}(\tau' - x) \right) \\ & + \psi \left(\frac{1}{2} + \frac{1}{y} + \frac{i}{\pi}(\tau' - x) \right) \\ & \left. - \psi \left(\frac{1}{2} + \frac{1}{y} + \frac{i}{\pi}(\tau' + x) \right) \right] \\ & + \frac{1}{\pi^2} \left[\psi' \left(\frac{1}{2} + \frac{1}{y} - \frac{i}{\pi}\tau' \right) \right. \\ & \left. - \psi' \left(\frac{1}{2} + \frac{1}{y} + \frac{i}{\pi}\tau' \right) \right]. \end{aligned}$$

Here, $\phi_0 = 2\pi^2\alpha/(\omega_0^2\beta^2)$, $x = \pi d/c\beta$, $y = \omega_c\beta$, $\tau' = \pi\tau/\beta$, $\psi(z)$ is the the digamma function, and $\psi'(z)$ its first derivative.

To proceed, we assume a high-temperature or strong-coupling regime, such that the dominant contributions to the integrals in Eqs. (A3) and (A4) will come from the peak in $\tilde{\phi}(\tau)$ at $\tau = 0$. More specifically, inspection of $C(x, y, \tau')$ reveals that for $y \gg 1$ (the scaling limit of large ω_c), we require $\phi_0 \gg 1$ for large x (weak correlations), or $\phi_0 x^2 \gg 1$ for small x (strong correlations), in order for an expansion of $\tilde{\phi}(\tau)$ around $\tau = 0$ to be valid. These definitions of the high-temperature (or strong-coupling) regime tally with the expansion parameters identified in Ref. 11. In the opposite limit, $y \ll 1$, we generally need $\phi_0 x^2 y^3 / \pi^4 \gg 1$, except in the limit of very large separations (vanishing correlations), $x \rightarrow \infty$, where $\phi_0 y \gg 1$ is the relevant condition.

With these conditions in mind, we therefore expand $\tilde{\phi}(\tau)$ to second order in $\tau' = \pi\tau/\beta$, which gives

$$\tilde{\phi}(\tau) \approx \phi_0 [C_0(x, y) - \tau'^2 C_2(x, y)], \quad (\text{A6})$$

where

$$\begin{aligned} C_0(x, y) = & \frac{i}{\pi x} \left[\psi \left(\frac{1}{2} + \frac{1}{y} + \frac{ix}{\pi} \right) - \psi \left(\frac{1}{2} + \frac{1}{y} - \frac{ix}{\pi} \right) \right] \\ & + \frac{2}{\pi^2} \psi' \left(\frac{1}{2} + \frac{1}{y} \right), \end{aligned} \quad (\text{A7})$$

and

$$C_2(x, y) = \frac{i}{2\pi^3 x} \left[\psi'' \left(\frac{1}{2} + \frac{1}{y} + \frac{ix}{\pi} \right) - \psi'' \left(\frac{1}{2} + \frac{1}{y} - \frac{ix}{\pi} \right) \right] + \frac{1}{\pi^4} \psi''' \left(\frac{1}{2} + \frac{1}{y} \right). \quad (\text{A8})$$

In this limit, Eqs. (A3) and (A4) are dominated by the terms containing a factor of $\exp[\tilde{\phi}(\tau)]$, allowing us to write

$$\gamma_{II}(\eta) \approx \frac{B^2 e^{\beta\eta/2} \beta}{2\pi} e^{\phi_0 C_0} \int_{-\infty}^{+\infty} d\tau' e^{i\tau'\eta\beta/\pi} e^{-\tau'^2 \phi_0 C_2}. \quad (\text{A9})$$

The integral is Gaussian and we arrive at the result

$$\gamma_{II}(\eta) = \frac{\beta B^2 e^{\phi_0 C_0(x, y)}}{2\sqrt{\pi} C_2(x, y) \phi_0} e^{\beta\eta/2} e^{-\beta^2 \eta^2 / [4\pi^2 C_2(x, y) \phi_0]}, \quad (\text{A10})$$

which reduces to

$$\gamma_{II}(\eta) \approx \gamma_{II}(0) = \frac{\beta B^2 e^{\phi_0 C_0(x, y)}}{2\sqrt{\pi} C_2(x, y) \phi_0}, \quad (\text{A11})$$

if the temperature is high enough such that $1/\eta\beta \gg 1$.

It remains now to determine the bath renormalization factor B , given by Eq. (7). To do so it is helpful to separate vacuum and thermal contributions. We write $B = B_0 B_{\text{th}}$, where

$$B_0 = \exp \left\{ - \int_0^\infty d\omega \frac{J(\omega)}{\omega^2} [1 - F_D(\omega, d)] \right\}, \quad (\text{A12})$$

and

$$B_{\text{th}} = \exp \left\{ - \int_0^\infty d\omega \frac{J(\omega)}{\omega^2} [1 - F_D(\omega, d)] [\coth(\beta\omega/2) - 1] \right\}. \quad (\text{A13})$$

Inserting the spectral density, and again assuming system-bath coupling in three dimensions, we find

$$B_0 = \exp \left[-\alpha \frac{\omega_c^2}{\omega_0^2} \left(\frac{(d\omega_c/c)^2}{1 + (d\omega_c/c)^2} \right) \right], \quad (\text{A14})$$

and

$$B_{\text{th}} = \exp \left[\frac{\phi_0}{2\pi^2} \left(\frac{i\pi}{x} [H(y^{-1} - ix/\pi) - H(y^{-1} + ix/\pi)] - 2\psi'(1 + y^{-1}) \right) \right], \quad (\text{A15})$$

where $H(m) = \sum_{i=1}^m (1/i)$ is the m th harmonic number. We note that in the infinite separation (uncorrelated) limit, one finds $B_0(d \rightarrow \infty) = \exp[-\alpha(\omega_c^2/\omega_0^2)]$, and

$$B_{\text{th}} = \exp \left(\alpha (\omega_c^2/\omega_0^2) \{1 - y^{-2} [\zeta(2, 1 + y^{-1}) + \zeta(2, y^{-1})]\} \right), \quad (\text{A16})$$

where $\zeta(s, a)$ is the generalized Riemann zeta function.

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