

Shortcuts to adiabaticity with a quantum control field

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Quantum adiabatic dynamics is the crucial element of adiabatic quantum computing and quantum annealing. Shortcuts to adiabaticity enable acceleration of the computational time by suppressing unwanted non-adiabatic processes with designed classical fields. Here, we consider quantum state transfer in the Landau-Zener model, which exemplifies the key elements of quantum adiabatic dynamics. We argue that non-adiabatic transitions can be suppressed by autonomous quantum dynamics, which involves coupling the Landau-Zener qubit to a second quantum system. By tuning the coupling strength, the composite quantum dynamics can reduce the probability of unwanted processes by more than two orders of magnitude. This is a prime example of control where the quantum properties of the control fields are key for implementing shortcuts to adiabaticity.

Introduction.—The Landau-Zener (LZ) model is a paradigm of a time-dependent quantum mechanical process [1–3] that is exactly solvable and encompasses the fundamental principles of quantum adiabatic transfer. For these reasons, the LZ model provides key insights into quantum thermodynamic principles [4], quantum phase transitions [5, 6], and quantum annealing dynamics [7–9], serving as the workhorse for identifying strategies for accelerating adiabatic transfers [5] and unveiling the mechanisms that determine their ultimate speed limits [10, 11]. Some of these strategies utilize classical control fields, which induce processes that suppress adiabatic transitions, thereby implementing so-called “shortcut to adiabaticity” protocols [12, 13]. This can be enforced through nonlinearities [14–17]. Perfect cancellation of the non-adiabatic transitions is achieved by counterdiabatic protocols [18–23], among which the control field is known for the LZ dynamics [19].

Thermal reservoirs can be beneficial for adiabaticity [24–27]. This observation has led to the identification of open-quantum-system approaches to shortcuts to adiabaticity [28–32] and to the formulation of strategies for reservoir-induced transitionless quantum driving [33, 34]. In these protocols, the incoherent dynamics is engineered through the coupling with a second quantum system, which can possess the features of a reservoir [35, 36] or undergo dissipative dynamics itself [32, 37–39]. The majority of these models disregard quantum correlations between the LZ qubit and the second quantum system, treating the latter as memoryless. Nevertheless, there are indications that memory effects can lead to more favorable speed limits [34, 39–41]. In particular, it has been posited that entanglement between the LZ qubit and the *quantum* control field during the dynamics can be resourceful for shortcuts to adiabaticity [39]. This hypothesis is corroborated by studies of the multilevel LZ model [30, 42, 43] and on quantum annealing protocols designed by auxiliary systems [44, 45].

In this work, we analyze the fidelity of quantum state transfer in the LZ model. We assume that the LZ qubit

couples linearly to a second quantum system, which we interchangeably denote by *quantum field* or *spectator*. We demonstrate that, when the LZ qubit dynamics is deep in the diabatic regime, where quantum state transfer is inefficient, the composite dynamics of the LZ qubit and spectator can correct for the diabatic processes and reduce the probability that they occur by more than two orders of magnitude. The protocol is illustrated in Fig. 1. Superadiabatic transfer is realized when the coupling between qubit and spectator is in the so-called ultra-strong coupling regime [46, 47]. The protocol does not require the time-dependent control of the quantum field; it is optimized by choosing the field eigenfrequency and the coupling strength with the LZ qubit, and is robust against fluctuations in the physical parameters.

The model.—The dynamics of the LZ qubit and field is governed by the Hamiltonian acting on the composite Hilbert space of qubit (s) and field (f):

$$\hat{H} = \hat{H}_s(t) \otimes \hat{\mathbb{I}}_f + \hat{H}_{\text{int}} + \hat{\mathbb{I}}_s \otimes \hat{H}_f, \quad (1)$$

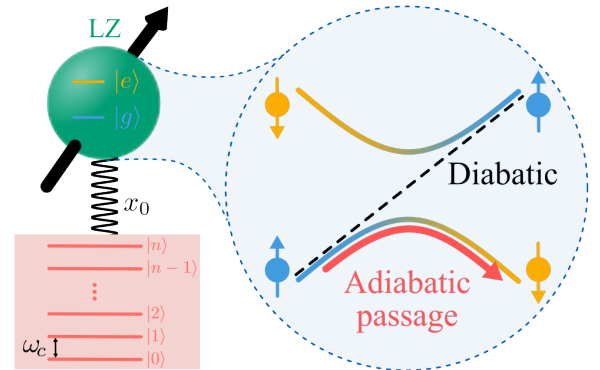


FIG. 1. Interference-assisted superadiabaticity: The ultra-strong coupling of an LZ qubit with a quantum field (bosonic or fermionic) leads to high-fidelity adiabatic passage even for fast sweeps, where the isolated LZ dynamics would be in the diabatic regime. Here, x_0 scales the strength of the coupling, and ω_c is the eigenfrequency of the quantum field.

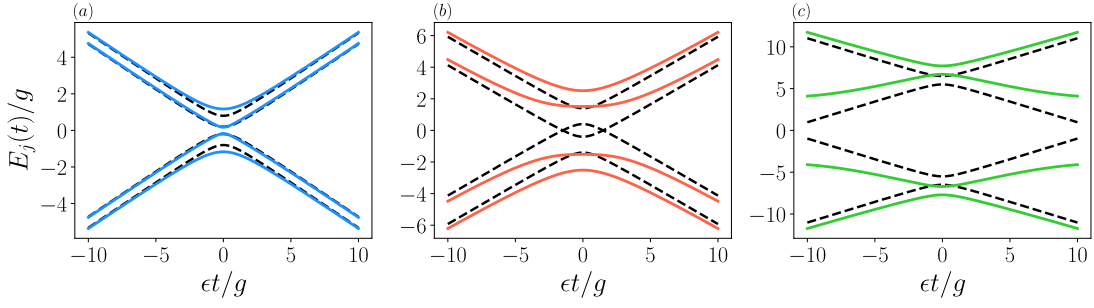


FIG. 2. Instantaneous energy branches $E_j(t)$ of the qubit-field Hamiltonian \hat{H} (1) (solid lines) for exemplary values of x_0 and ω_c corresponding to three regimes: (a) Regime (I) $x_0 = \omega_c = 0.6g$, where the coupling with the field tends to close the minimal gap; (b) Regime (II) $x_0 = \omega_c = 1.8g$, where the minimal gap increases with Δ and the eigenenergies flatten in the vicinity of $t = 0$; (c) Regime (III), for $x_0 = 4g$ and $\omega_c = 12g$, gaps start to close at $|t| > 0$. In all cases, $\epsilon = 2g^2$. The dashed lines are the instantaneous eigenenergies of the qubit-field system at $x_0 = 0$.

where $\hat{\mathbb{I}}_i$ is the identity operator in the corresponding Hilbert space, see Fig. 1. The LZ Hamiltonian of the qubit reads

$$\hat{H}_s(t) = \frac{1}{2} [\epsilon t \hat{\sigma}^z + g \hat{\sigma}^x], \quad (2)$$

where $\hat{\sigma}^j$ are the Pauli matrices ($j = x, y, z$). Hamiltonian (2) describes the dynamics of a spin 1/2 (qubit) in the presence of a classical magnetic field with a constant component along x and a time-dependent amplitude along z , which varies at a constant (sweep) rate ϵ . The instantaneous Hamiltonian $\hat{H}_s(t)$ is diagonal in the eigenbasis $|\pm\rangle_t$ with eigenenergies $E_{\pm}(t) = \pm \frac{1}{2} \sqrt{\epsilon^2 t^2 + g^2}$. In the absence of the spectator field ($x_0 = 0$), the probability that the qubit is excited into the upper branch $|+\rangle_t$ is the infidelity $\mathcal{I}_{LZ} = |\langle + | \psi \rangle_t|^2$, where the quantum state $|\psi\rangle_t$ solves the time-dependent Schrödinger equation $i\partial_t |\psi\rangle_t = \hat{H}(t) |\psi\rangle_t$ (we set $\hbar = 1$). For $t \rightarrow \infty$, the infidelity takes the compact form $\mathcal{I}_{LZ} = \exp(-\pi g^2/2\epsilon)$, showing that adiabatic transfer is warranted for $g^2/\epsilon \gg 1$. Evidently, \mathcal{I}_{LZ} decreases monotonously by increasing the gap g and/or decreasing the sweep rate ϵ [42].

In what follows, we consider the diabatic regime, in which $g^2/\epsilon < 1$ and $\mathcal{I}_{LZ} > 0.2$. For convenience, we now assume that the spectator field is a second qubit at the transition frequency ω_c , with $\hat{H}_f = \omega_c \hat{\tau}^z/2$, and $\hat{\tau}^j$ denoting the Pauli matrices acting on the states of the spectator Hilbert space. The qubit-spectator interaction is the minimal coupling Hamiltonian

$$\hat{H}_{\text{int}} = x_0 \hat{\sigma}^x \otimes \hat{\tau}^x, \quad (3)$$

and is scaled by the parameter x_0 with the dimension of an energy.

Hamiltonian (1), and variations thereof, have been discussed in several contexts. A trivial but instructive case is found for $\epsilon = 0$. Then, the interaction Hamiltonian commutes with the qubit Hamiltonian, and the spectator field undergoes a dynamics conditioned by the state

of the qubit. Measurement of the state of the spectator then provides information on the state of the qubit. This is the basic principle underlying quantum non-demolition (QND) measurements [48–50]. This concept has been extended in Refs. [32, 39] for designing adiabatic protocols in which the instantaneous interaction commutes with $\hat{H}_s(t)$ at all times. For $\epsilon \neq 0$, when the spectator's eigenfrequency vanishes $\omega_c = 0$, the spectator field dynamics is stationary in the eigenbasis of $\hat{\tau}^x$. By judiciously selecting the state of the spectator field, the qubit undergoes a LZ dynamics with the effective gap $g' = g + x_0$ at the anticrossing $t = 0$. This results in a rescaling of the adiabaticity parameter g'^2/ϵ , which increases the fidelity of the transfer (see also [16] for a similar case). In its complete form, the dynamics governed by Hamiltonian (1) spans the four-dimensional Hilbert space of qubit and spectator. In what follows, we show that this is key for promoting superadiabatic quantum state transfer.

Adiabatic basis and adiabatic theorem.—Relevant frequency scales of the dynamics governed by the Hamiltonian in Eq. (1) can be extracted in the adiabatic basis of the instantaneous eigenstates of $\hat{H}(t)$. Figure 2 displays the eigenenergies as a function of t for different values of x_0 and ω_c . In general, we observe a doubling of the two branches $E_{\pm}(t)$ of the LZ model. The coupling x_0 entangles the qubit and spectator states. At fixed x_0 , tuning ω_c can give rise to a single anticrossing at $t = 0$ or to multiple anticrossings at $t = 0$ and $|t| > 0$.

Some insight is gained in the two limiting cases: $|t| \rightarrow \infty$ and $t = 0$. In the trivial case $|t| \rightarrow \infty$, the qubit and spectator are decoupled and the Hamiltonian is diagonal in the basis of eigenstates of $\hat{\sigma}^z \otimes \hat{\tau}^z$. At $t = 0$, the Hamiltonian can be exactly diagonalized. In fact, the instantaneous Hamiltonian commutes with $\hat{\sigma}^x$ and the qubit eigenstates are the instantaneous states $|\pm\rangle_0$. Each branch is split into two values: For $|+\rangle_0$, the eigenvalues are $E_{\pm}^{(+)} = +g/2 \pm \Delta/2$, while for $|-\rangle_0$ they read $E_{\pm}^{(-)} = -g/2 \pm \Delta/2$. Thus, the two key quantities determining the energy splitting at $t = 0$ are the minimum gap g of the

LZ Hamiltonian (2) and the spectator splitting energy

$$\Delta = \sqrt{4x_0^2 + \omega_c^2}. \quad (4)$$

This simple analysis identifies a threshold value $\Delta = \Delta_c^{(1)} \approx g$, at which the spectrum becomes doubly degenerate. This threshold value separates two different regimes: For $\Delta < \Delta_c^{(1)}$ the coupling with the spectator field tends to close the minimal gap between $|\pm\rangle_0$, as visible in Fig. 2(a). In this regime, the gap is $\sim |g - \Delta|$, and the efficiency of the adiabatic protocol worsens with increasing Δ . This tendency inverts for $\Delta > \Delta_c^{(1)}$. In this regime, the coupling with the spectator field increases the energy splitting at $t = 0$ between the states $|+\rangle_0$ and $|-\rangle_0$. Inspection of Fig. 2(c) shows that, for larger values of x_0 , the gap between upper and lower branches closes at $t \neq 0$. Considerations on the structure of the adiabatic eigenvalues show that this occurs at $|t| = t_\epsilon \sim \Delta/\epsilon$ when $\Delta \gtrsim \epsilon^2$. This identifies a second threshold value $\Delta_c^{(2)}(\epsilon)$, above which the coupling with the quantum field shall promote diabatic transitions. Interestingly, the threshold $\Delta_c^{(2)}(\epsilon)$ increases monotonically with the sweep rate ϵ . In the diabatic regime, for $g < \sqrt{\epsilon}$, this second threshold is then an upper bound: $\Delta_c^{(2)}(\epsilon) > \Delta_c^{(1)}$, and there exists a finite range of values x_0 and ω_c for which $\Delta \in [\Delta_c^{(1)}, \Delta_c^{(2)}]$.

This led us to identify three regimes induced by the coupling with the quantum field, which are strictly valid for $g > \sqrt{\epsilon}$. The regime (I) where $0 < \Delta < \Delta_c^{(1)}$. Here, the minimal energy gap decreases monotonically by increasing Δ and the diabatic transitions occur with the largest probability at $t \sim 0$. Regime (II) is found at $\Delta \in (\Delta_c^{(1)}, \Delta_c^{(2)})$. Here, increasing Δ increases the minimal gap and is thus expected to suppress diabatic transitions. In regime (III), $\Delta > \Delta_c^{(2)}$, diabatic transitions might also occur at $|t| = t_\epsilon > 0$. Since the bound $\Delta_c^{(2)}$ increases with ϵ , entering regime (III) necessitates sufficiently large values of Δ for the condition $\Delta > \Delta_c^{(2)}$ to be satisfied. Thus, increasing the sweep rate ϵ is expected to increase the size of the superadiabatic region (II) in parameter space. This consideration might seem contradictory. In fact, diabatic effects scale with $\partial_t \hat{H}(t)$ and thus with the sweep rate ϵ . According to the adiabatic theorem, adiabatic transfer is warranted when

$$\frac{|\langle i(t) | \partial_t \hat{H}(t) | j(t) \rangle|}{(E_i(t) - E_j(t))^2} \ll 1,$$

with i, j labeling the instantaneous eigenstates and $E_{i,j}$ the corresponding eigenenergies. For $x_0 = 0$, this reduces to the inequality $\epsilon/g^2 \ll 1$. By coupling with the spectator field, the numerator becomes generally smaller because of the reduced overlap between each pair of states. Moreover, in regime (II), the denominator becomes larger. We therefore expect that the fidelity of the

transfer will increase when the parameters of the spectator qubit are chosen in this regime.

Dynamics.— We use the following figures of merit to characterize the dynamics. Given $|\Psi\rangle_t$, the state of the composite system of qubit and spectator at time t , we determine the density matrix of the qubit $\hat{\rho}(t) = \text{Tr}_f\{|\Psi\rangle_t\langle\Psi|\}$ by tracing out the degrees of freedom of the spectator field. We then define the infidelity at time t as the probability $P(t)$ that the qubit is in the excited branch $|+\rangle_t$ of the bare LZ Hamiltonian:

$$P(t) = \text{Tr}\{\hat{\rho}(t)|+\rangle_t\langle+\rangle_t\}. \quad (5)$$

The probability $P(t)$ refers solely to the excitation of the LZ model, irrespective of the state of the spectator field. Entanglement between qubit and spectator is assessed by means of the purity

$$\gamma(t) = \text{Tr}\{\hat{\rho}^2(t)\}. \quad (6)$$

The purity is a witness of entanglement via its link with the Rényi entropy $S_2 = -\ln \gamma$; its maxima $\gamma = 1$ signal that the qubit is in a pure state.

Figure 3 displays the evolution of $P(t)$ for $\epsilon = 2g^2$. For this choice, and in the absence of the spectator field ($x_0 = 0$), the probability that the qubit is excited into the upper branch at the end of the transfer exceeds 0.5. The coupling with the spectator for a choice of parameters in regime (II) leads to a dramatic reduction in the transition probability, performing oscillations which are damped out for larger times (not shown). Correspondingly, the purity reaches values close to unity. During the protocol, the purity is visibly reduced at the anticrossing ($t = 0$), showing that entanglement between the qubit

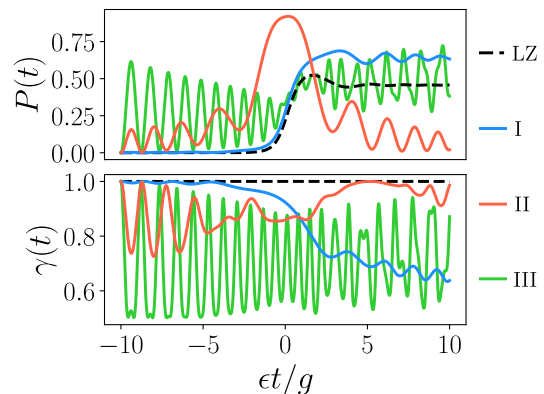


FIG. 3. Evolution of the transition probability $P(t)$ (5) and the purity $\gamma(t)$ (6) as a function of time for the qubit-field dynamics governed by Hamiltonian (2) and $\epsilon = 2g^2$. Three different choices of x_0 and ω_c are shown, illustrating the regimes of the annealing dynamics: (I) $x_0 = g/4$ and $\omega_c = 2x_0$, (II) $x_0 = 2g$ and $\omega_c = x_0/4$, (III) $x_0 = 4g$ and $\omega_c = 4x_0$. The black dashed line represents the Landau-Zener case ($x_0 = 0$). The qubit-field system is initialized at time $t_i = -10g/\epsilon$ in state $|\psi_0\rangle = |-\rangle_{t_i} \otimes |\downarrow\rangle$ and evolved up until time $t_f = -t_i$.

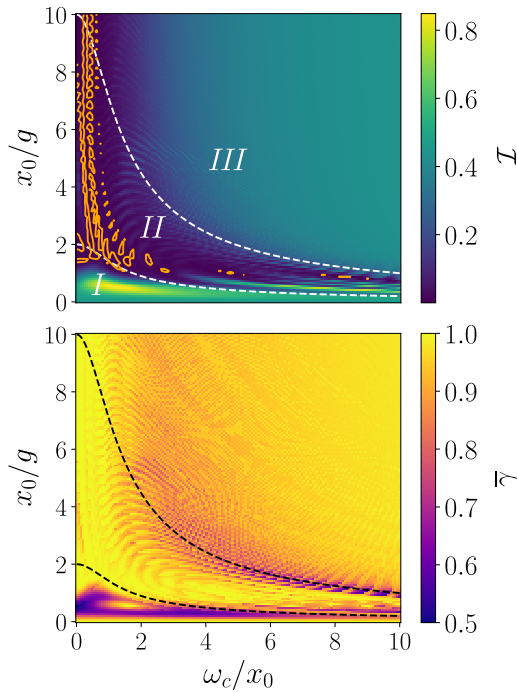


FIG. 4. Adiabatic transfer induced by a quantum field: Infidelity \mathcal{I} and purity $\bar{\gamma}$ of the final state as a function of the parameters x_0/g and x_0/ω_c . Results are obtained for the same annealing procedure and choice of parameters as for Fig. 3. Other parameters are fixed such that $g^2/\epsilon = 0.5$. The duration of the annealing protocol is optimized such that $P(t_f)$ is at the nearest local minimum to $t = 10g/\epsilon$. The dashed lines indicate the critical values $\Delta_c^{(1)} \approx g$ and $\Delta_c^{(2)}$, separating the three regimes (I), (II), and (III). The orange contour delimits regions where the infidelity lies below 5×10^{-3} .

and spectator is instrumental in achieving high fidelity at the end of the transfer.

Figure 4 visualizes the central result of this work. It displays the infidelity and purity at $t_f \approx 10g/\epsilon$ and as a function of x_0 (in units of g) and of ω_c/x_0 . We emphasize that we choose a relatively short transfer time, $t_f \approx 10g/\epsilon$, where the infidelity still oscillates. The choice of $\epsilon = 2g^2$ fixes the asymptotic value of the infidelity of the LZ transfer above 0.5. The dashed lines indicate the position of the two thresholds $\Delta_c^{(1)}$ and $\Delta_c^{(2)}$. As expected from the analysis of the spectrum in the adiabatic basis, the transfer is inefficient in regime (I) at $\Delta < \Delta_c^{(1)}$. In contrast, in regime (II), the infidelity drops to values below 0.05. The bright contours highlight the regions where it reaches below 10^{-3} : This occurs in islands distributed along an optimal value of Δ as well as along a stripe at $\omega_c/x_0 \lesssim 1$. Regime (III) is generally characterized by low fidelity. An exception is the stripe at $\omega_c/x_0 \lesssim 1$, which stretches into this region. Here, $x_0 \gg g$. For $\omega_c = 0$, the effect of the spectator is simply to renormalize the minimal gap to $g' = g + x_0 \approx x_0$. As a consequence, even though in the absence of the coupling

the sweep is in the diabatic regime ($\epsilon/g^2 > 1$), adding the field effectively yields $\epsilon/g'^2 \ll 1$ and adiabaticity is re-established. The minimal infidelity, however, is reached for ω_c of the order of x_0 , indicating that the quantum nature of the spectator is essential to achieve high transfer efficiencies. We note that the initial state of the spectator is not important: Choosing different initial values only affects the infidelity at $\omega_c = 0$, while for $\omega_c > 0$ the diagram is qualitatively the same.

We have varied the form of the coupling of the qubit in Eq. (3) and verified that the superadiabatic transfer occurs independently of the choice of the Pauli matrix $\hat{\sigma}_x$ or $\hat{\sigma}_y$. This indicates that the superadiabaticity we report emerges from a very different dynamics than the counteradiabatic protocol of Ref. [19], which instead strictly requires a classical control field along $\hat{\sigma}_y$ for perfect cancellation. In our case, superadiabaticity is due to interference between the different paths corresponding to the branches in the adiabatic basis, as visible in the evolution of the purity, Fig. 3, and by the corresponding reduction of the amplitude of the diabatic transitions. We have additionally determined the transfer infidelity when the spectator is a harmonic oscillator at frequency ω_c , taking the interaction Hamiltonian $\hat{H}_{\text{int}} = x_0 \hat{\sigma}^x \otimes (\hat{a} + \hat{a}^\dagger)$, with \hat{a} , \hat{a}^\dagger the annihilation and creation operators. We observe superadiabatic behavior qualitatively similar to Fig. 4. Based on these insights, we denote by *interference-assisted superadiabaticity* (IAS) the superadiabatic transfer induced by the quantum field, emphasizing the role of quantum coherence and entanglement in this dynamics.

Discussion.— The dynamics discussed here could be realized in solid-state platforms [43, 51, 52], which offer the possibility to selectively couple a qubit with a spectator field. These versatile platforms allow for coupling strengths in the ultra-strong coupling regime, meeting the criteria to be in the optimal IAS regime. Our results show that fluctuations of 10% of the parameter values do not significantly affect the efficiency (see Fig. 4), demonstrating that the protocol is robust against variation of the experimental parameters. We have verified that moderate dissipation of the spectator leads to a small reduction in the efficiency, but can be advantageous as it “damps down” the oscillations after the transfer.

A striking feature of IAS is the protocol’s remarkable efficiency deep in the diabatic regime: Almost unit fidelity can be reached at fast sweeps. The dynamics is reminiscent of counteradiabatic protocols, even though we have no evidence of transitionless driving. The infidelity could be further reduced by optimizing the schedule for the qubit sweep, e.g. by implementing local adiabatic evolution, extending the approach of Ref. [53] to the coupled qubit-spectator system. Scalability to a qubit string requires the capability to implement ultra-strong global coupling between qubits and the spectator. This can be done by means of a central spin or oscillator

model as realized in CQED platforms [54]. The robustness against fluctuations of the control field relaxes the requirements on the precision about the knowledge of the annealing gap. Our study is in line with the insights of recent works [44, 45, 55, 56] and points toward the advantage of autonomous quantum protocols for quantum supremacy.

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