Remote Charging and Degradation Suppression for the Quantum Battery

Wan-Lu Song,[1] Hai-Bin Liu,[1] Bin Zhou,[1,*] Wan-Li Yang,[2,†] and Jun-Hong An[3,‡]

[1] *Department of Physics, Hubei University, Wuhan 430062, China*

[2] *State Key Laboratory of Magnetic Resonance and Atomic and Molecular Physics, Innovation Academy for Precision Measurement Science and Technology, Chinese Academy of Sciences, Wuhan 430071, China*

[3] *Key Laboratory of Quantum Theory and Applications of MoE, Lanzhou Center for Theoretical Physics, and Key Laboratory of Theoretical Physics of Gansu Province, Lanzhou University, Lanzhou 730000, China*

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The quantum battery (QB) makes use of quantum effects to store and supply energy, which may outperform its classical counterpart. However, there are two challenges in this field. One is that the environment-induced decoherence causes the energy loss and aging of the QB, the other is that the decreasing of the charger-QB coupling strength with increasing their distance makes the charging of the QB become inefficient. Here, we propose a QB scheme to realize a remote charging via coupling the QB and the charger to a rectangular hollow metal waveguide. It is found that an ideal charging is realized as long as two bound states are formed in the energy spectrum of the total system consisting of the QB, the charger, and the electromagnetic environment in the waveguide. Using the constructive role of the decoherence, our QB is immune to the aging. Additionally, without resorting to the direct charger-QB interaction, our scheme works in a way of long-range and wireless-like charging. Effectively overcoming the two challenges, our result supplies an insightful guideline to the practical realization of the QB by reservoir engineering.

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Introduction.—The miniaturization of electronic devices has inspired a rapid development of quantum technologies, which are expected to bring a lot of technological innovations [1–3]. Among them, quantum thermodynamics [4–6] has emerged as a field aiming to reconstruct thermodynamics via basic laws of quantum mechanics and to break through the constraint of classical physics to device performance by quantum effects [7–12]. As an energy-storing and -converting device in atomic size, the quantum battery (QB) promotes the application of quantum effects in thermodynamics [13,14]. With the aid of quantum resources, such as quantum coherence and entanglement, the QB possesses a potential of stronger charging power, higher charging capacity, and larger work extraction than its classical counterpart [15–21].

An efficient QB scheme is a prerequisite for its realization and further applications. The widely studied QB model is based on two-level systems (TLSs), which are charged by the coherent coupling to other TLSs [22–26] or external fields [27–33]. Experiments on the single-body QB were carried out in nuclear magnetic resonance [34], superconducting circuit [35,36], semiconducting quantum dot [37,38], and photonic [39] systems. It was also found that the collection of more TLSs exhibits more energy storage capacity and the collective charging benefiting from quantum correlation induces stronger charging power of the QB [40–61]. However, the practical performance of the QB is challenged by two facts. One is that its energy during the long-time storage process is depleted by the environment-induced decoherence, which is called the aging of the QB [62]. The other is that its charging becomes inefficient when the charger-QB coupling is weakened due to certain reasons. Several schemes using feedback control [63], dark states [64], Floquet engineering [65], and environment engineering [66–77], have been proposed to beat the decoherence effects. However, how to suppress the decoherence effects on the QB and improve the charging efficiency simultaneously is still an open question.

Inspired by the development of quantum interconnect to establish correlation among distant quantum entities [78], we propose a remote-charging QB scheme via coupling the QB and the charger to a rectangular hollow metal waveguide [79–81]. Utilizing the non-Markovian decoherence induced by the electromagnetic field (EMF) in the waveguide, we find a mechanism that allows for a persistent energy exchange between the noninteracting charger and QB in the steady state. It is due to the formation of two bound states in the energy spectrum of the total system formed by the charger, the QB, and the EMF. Not causing the aging of the QB, the decoherence in our scheme plays a constructive role in the charging. Without resorting to the direct charger-QB coupling, our charging works in a long-range and wireless way. Overcoming the aging and low charging efficiency in the weak charger-QB coupling

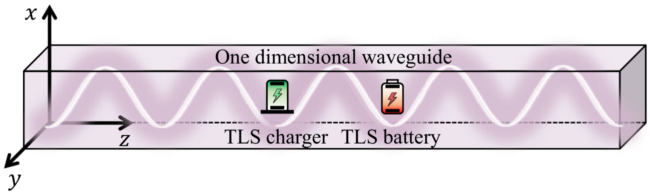

FIG. 1. Scheme of a remote-charging QB. The noninteracting charger and QB formed by TLSs are separated in a distance Δ_z in a rectangular hollow metal waveguide. An efficient remote charging to the QB is realized by the mediation role of the EMF in the waveguide.

condition simultaneously, our scheme supplies an insightful guideline to the practical realization of the QB.

Remote-charging QB model.—We consider a long-range charging QB scheme, which consists of two TLSs acting as the charger and the QB residing in a rectangular hollow metal waveguide (see Fig. 1). Without direct interaction, the charger and QB couple to a common EMF in the waveguide. The Hamiltonian reads [80–82]

$$\hat{H} = \omega_0 \sum_{j=1,2} \hat{\sigma}_j^\dagger \hat{\sigma}_j + \sum_{k=-\infty}^{+\infty} \Big[\omega_k \hat{a}_k^\dagger \hat{a}_k + \sum_{j=1,2} (g_{k,j} \hat{\sigma}_j^\dagger \hat{a}_k + \text{H.c.}) \Big], \tag{1}$$

where $\hat{\sigma}_j = |g_j\rangle\langle e_j|$ is the transition operator from the excited state $|e_j\rangle$ to the ground state $|g_j\rangle$ with frequency ω_0 of the charger when $j=1$ and of the QB when $j=2$ and $\hat{a}_k$ is the annihilation operator of the kth mode with frequency ω_k of the EMF. The coupling strength between the jth TLS and the kth mode of the EMF is $g_{k,j} = \sqrt{\omega_k/(2\epsilon_0)}\mathbf{d}_j \cdot \mathbf{u}_k(\mathbf{r}_j)$, where ϵ_0 is the vacuum permittivity, $\mathbf{d}_j$ is the dipole moment of the jth TLS, and $\mathbf{u}_k(\mathbf{r}_j)$ is the spatial function of the EMF. The dispersion relation of the EMF is $\omega_k = [(ck)^2 + \omega_{mn}^2]^{1/2}$, where c is the speed of light, k is the longitudinal wave number, and $\omega_{mn} = c[(m\pi/a)^2 + (n\pi/b)^2]^{1/2}$, with a and b being the transverse lengths of the waveguide, is the cutoff frequency of the transverse mode of the EMF. Assuming that the two TLSs are polarized in the z direction, i.e., $\mathbf{d}_1 = \mathbf{d}_2 = d_z\mathbf{z}$, and considering only the dominated EMF mode with $m=n=1$, the coupling strength is further characterized by the spectral density

$$J_{|j-j'|}(\omega) = \frac{\Gamma_{11}}{2\pi} \frac{\cos\left[\frac{z_j - z_{j'}}{c}\sqrt{\omega^2 - \omega_{11}^2}\right]}{\sqrt{(\omega/\omega_{11})^2 - 1}} \Theta(\omega - \omega_{11}). \tag{2}$$

Here, $\Gamma_{11} = (4\omega_{11}d_z^2/\epsilon_0 abc)\sin^2(\pi x_0/a)\sin^2(\pi y_0/b)$ is the radiation rate of the TLS with transverse position (x_0, y_0) into the waveguide and $\Theta(\omega - \omega_{11})$ is the Heaviside step function. The spectral density $J_0(\omega)$ induces an individual spontaneous emission and Lamb frequency shift to the jth TLS, while $J_1(\omega)$ induces a correlated spontaneous emission and dipole-dipole interaction between the two TLSs. Different from Refs. [24,47,54,59,72], we use such an incoherent correlation to charge the QB.

Our scheme begins from a pump laser to initialize the charger to its excited state. When a charging to the QB is needed, the QB is taken to the waveguide not necessarily too near to the charger. The EMF in the waveguide is expected to mediate a charger-QB energy transfer. This process under the condition $|\Psi(0)\rangle = |e_1, g_2; \{0_k\}\rangle$ is described by $|\Psi(t)\rangle = [\sum_j c_j(t)\hat{\sigma}_j^\dagger + \sum_k d_k(t)\hat{a}_k^\dagger]|g_1, g_2, \{0_k\}\rangle$. Here, $c_j(t)$ and $d_k(t)$ are the excited-state probability amplitudes of the TLSs and the EMF with only one photon in the kth mode, respectively. From the Schrödinger equation, we obtain [83]

$$\dot{c}_j(t) + i\omega_0 c_j(t) + \sum_{j'=1,2} \int_0^t f_{|j-j'|}(t-\tau) c_{j'}(\tau) d\tau = 0, \tag{3}$$

where $f_{|j-j'|}(x) = \int_0^\infty J_{|j-j'|}(\omega) e^{-i\omega x} d\omega$, $c_1(0) = 1$, and $c_2(0) = 0$. Equation (3) reveals that, although the direct charger-QB interaction is absent, an effective charger-QB coupling is induced by the mediation role of the EMF in the waveguide. The convolution in Eq. (3) renders the dynamics non-Markovian. The energy of the QB is

$$\mathcal{E}(t) = \text{Tr}[\rho_B(t)\hat{H}_B], \tag{4}$$

where $\hat{H}_B = \omega_0 \hat{\sigma}_2^\dagger \hat{\sigma}_2$ and $\rho_B(t) = \text{Tr}_1[|\Psi(t)\rangle\langle\Psi(t)|]$. According to the second law of thermodynamics, not all of $\mathcal{E}(t)$ can be converted into work [85]. The maximum of the extractable energy called ergotropy is defined as [18,19]

$$\mathcal{W}(t) = \text{Tr}[\rho_B(t)\hat{H}_B] - \text{Tr}[\tilde{\rho}_B(t)\hat{H}_B], \tag{5}$$

where $\tilde{\rho}_B(t) = \sum_k r_k(t)|\varepsilon_k\rangle\langle\varepsilon_k|$ is the passive state, $r_k(t)$ are the eigenvalues of $\rho_B(t)$ ordered in a descending sort, and $|\varepsilon_k\rangle$ are the eigenstates of $\hat{H}_B$ with the corresponding eigenvalues ε_k ordered in an ascending sort.

In the special case of the weak TLS-EMF coupling, we apply the Markovian approximation to neglect the memory effect and extend the upper limit of the τ integral in Eq. (3) to infinity. The approximate solution is

$$c_j^{\text{MA}}(t) = e^{-(i\omega_0 + \Upsilon_0)t}[e^{-\Upsilon_1 t} - (-1)^j e^{\Upsilon_1 t}]/2, \tag{6}$$

where $\Upsilon_j = \pi J_j(\omega_0) + i\delta_j$, with $\delta_j = \mathcal{P}\int_0^\infty d\omega[J_j(\omega)/(\omega_0 - \omega)]$ and $\mathcal{P}$ denoting the Cauchy principal value. Equation (6) reveals that the EMF in the waveguide indeed can mediate an energy exchange between the charger and the QB. However, it is transient and the energies in both of them exponentially damp to zero in the long-time condition due to the destructive role of the Markovian decoherence.

In contrast to the above approximate result, the mediated charger-QB coupling can induce a persistently reversible energy exchange between them even in the steady state in the non-Markovian dynamics. Equation (3) is only numerically solvable. However, its long-time form is analytically obtainable by the Laplace transform method. It converts Eq. (3) to $\tilde{c}_j(s) = \{1/[\Xi(s) + \tilde{f}_1(s)] - (-1)^j/[\Xi(s) - \tilde{f}_1(s)]\}/2$, where $\Xi(s) = s + i\omega_0 + \tilde{f}_0(s)$ and $\tilde{f}_{|j-j'|}(s) = \int_0^\infty [J_{|j-j'|}(\omega)/(i\omega + s)]d\omega$. $c_j(t)$ is obtained by making the inverse Laplace transform to $\tilde{c}_j(s)$, which needs finding the poles of $\tilde{c}_j(s)$ from

$$Y_\pm(E) \equiv \omega_0 - \int_0^\infty \frac{J_0(\omega) \pm J_1(\omega)}{\omega - E} d\omega = E, \quad (E = is). \tag{7}$$

Note that the roots E of Eq. (7) are just the eigenenergies of Eq. (1). To prove this, we expand the eigenstate as $|\Phi\rangle = [\sum_j \alpha_j \hat{\sigma}_j^\dagger + \sum_k \beta_k \hat{a}_k^\dagger]|g_1, g_2, \{0_k\}\rangle$. From $\hat{H}|\Phi\rangle = E|\Phi\rangle$, we obtain $(E - \omega_0)\alpha_j = \sum_{j'} \int_0^\infty [J_{|j-j'|}(\omega)\alpha_{j'}/(E - \omega)]d\omega$, which readily results in Eq. (7) after eliminating α_j [83]. Thus, on one hand, Eq. (7) governs the dynamical evolution of $c_j(t)$, and on the other hand, it determines the eigenenergies of the total system. This means that the dynamics of the TLSs is essentially determined by the energy-spectrum feature of the total system. $Y_\pm(E)$ is ill-defined and jumps rapidly between $\pm\infty$ in the regime $E > \omega_{11}$ because of the divergence of the integrand. Thus, Eq. (7) has infinite roots in $E > \omega_{11}$, which form a continuous energy band. In the regime $E < \omega_{11}$, $Y_\pm(E)$ are monotonic decreasing functions of E and an isolated root E_+^b or E_-^b is formed provided $Y_+(\omega_{11}) < \omega_{11}$ or $Y_-(\omega_{11}) < \omega_{11}$. We call the eigenstates corresponding to $E_\pm^\mathrm{b}$ bound states. Substituting these poles from Eq. (7) into the inverse Laplace transform, the bound-state energies $E_\pm^\mathrm{b}$ contribute two nontrivial residues; while the band energies in the continuum contribute a branch cut, which tends to zero in the long-time limit due to the out-of-phase interference [86,87]. Thus, we have

$$c_j(\infty) = \begin{cases} 0, & 0 \text{ bound state,} \\ Z_+ e^{-iE_+^\mathrm{b} t}, & 1 \text{ bound state,} \\ Z_+ e^{-iE_+^\mathrm{b} t} - (-1)^j Z_- e^{-iE_-^\mathrm{b} t}, & 2 \text{ bound states,} \end{cases} \tag{8}$$

where $Z_\pm = \frac{1}{2}(1 + \int_0^\infty \{[J_0(\omega) \pm J_1(\omega)]/(E_\pm^\mathrm{b} - \omega)^2\} d\omega)^{-1}$ is the residue contributed by $E_\pm^\mathrm{b}$ to $c_j(t)$. Equation (8) indicates the decisive role played by the energy-spectrum feature of the total TLS-EMF system in the non-Markovian dynamics of the TLSs. It is remarkable to see that the formation of the bound states prevents $|c_j(t)|^2$ from damping to zero, which is not captured by the Markovian result in Eq. (6) and can be used to remotely charge the QB. When the QB is not needed for supplying energy to other devices, it is kept in the waveguide. Reversibly exchanging energy with the charger, the QB does not experience the aging caused by the energy loss in this storage process [83].

Equation (8) reveals that the QB is dynamically synchronous with the charger in the steady state. In the absence of the bound state, no energy is left in both the QB and the charger and the energy completely relaxes into the EMF of the waveguide. If one bound state with eigenenergy E_+^b is formed, then a stable energy $\mathcal{E}(\infty) = \omega_0 Z_+^2$ is preserved in the QB, which is equal to the one in the charger. As long as two bound states with eigenenergies $E_\pm^\mathrm{b}$ are formed, a periodic energy exchange of the QB with the charger in a frequency $E_+^\mathrm{b} - E_-^\mathrm{b}$, i.e., $\mathcal{E}(\infty) = \omega_0\{Z_+^2 + Z_-^2 - 2Z_+Z_- \cos[(E_+^\mathrm{b} - E_-^\mathrm{b})t]\}$, is kept in the steady state. Such a Rabi-like oscillation is naturally expected in the case of the direct charger-QB coupling. However, the charging realized by the direct coupling is generally fragile to the environment-induced decoherence, which results in the aging of the QB. Our result indicates that, different from the destructive role in the direct-coupling case, the non-Markovian decoherence caused by the common EMF can be used to realize an ideal charging to the QB in the open system framework. It is just the constructive role played by the decoherence that generates such an efficient charging in our scheme.

The persistent Rabi-like charger-QB energy exchange endows our scheme with a substantial difference from previous environment-assisted charging schemes either in the Markovian approximation [66–73] or in the non-Markovian dynamics [74–77]. Our charging performance, which is as perfect as the ideal charging case by direct charge-QB interaction, is guaranteed by the feature of the energy spectrum of the total system consisting of the charger, the QB, and the EMF in the waveguide. It makes our scheme immune to the environment-induced aging. Our scheme works in a way of long-range and wireless-like charging, while the spatial charger-QB distance has never been investigated. As an attractive feature in future applications, it overcomes the low-charging-efficiency problem experienced by the conventional QB schemes when the charger-QB coupling is weak.

Numerical results.—Choosing the charger-QB distance $\Delta_z \equiv |z_1 - z_2| = 0.1\lambda_{11}$, with $\lambda_{11} = 2\pi c/\omega_{11}$, which is much larger than the typical range of the dipole-dipole interactions, we plot in Fig. 2 the energy spectrum of the total system formed by the QB, the charger, and the EMF and the evolution of the QB energy $\mathcal{E}(t)$ in different frequency ω_0 of the TLSs. It is interesting to find that, in contrast to the damping to zero in the Markovian case, $\mathcal{E}(t)$ approaches a finite value for a moderate ω_0, while it exhibits a lossless Rabi-like oscillation for a small ω_0.

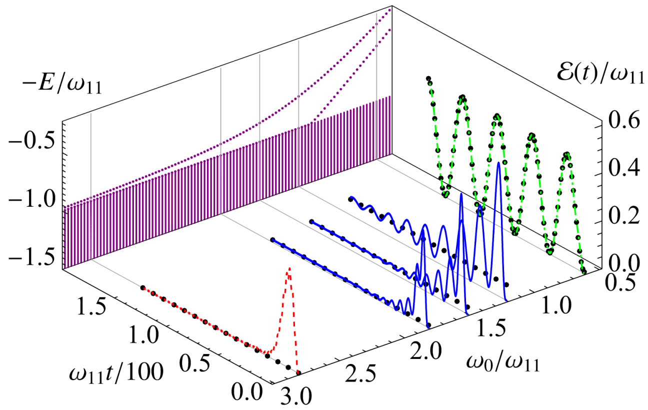

FIG. 2. Energy spectrum E of the total system via solving Eq. (7) and evolution of the QB energy $\mathcal{E}(t)$ via solving Eq. (3) in different ω_0 in the presence of zero (red dashed lines), one (blue solid lines), and two (green dot-dashed lines) bound states. The black dots mark the results evaluated from the analytical equation (8). We use $\Delta_z = 0.1\lambda_{11}$ and $\Gamma_{11} = 0.5\omega_{11}$.

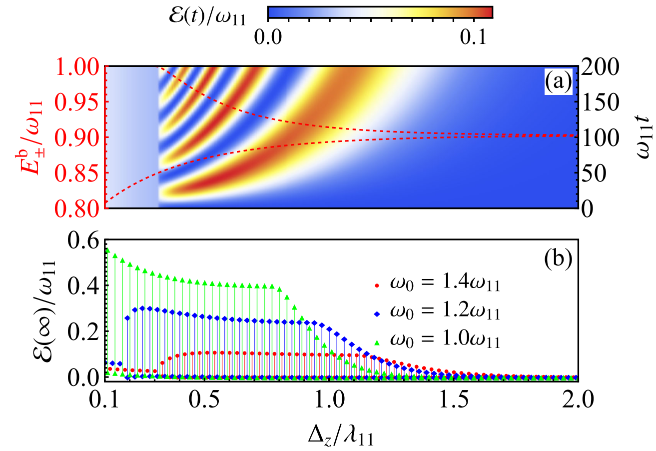

FIG. 3. (a) Bound-state eigenenergies $E^{\rm b}_{\pm}$ (red dashed lines) and long-time evolution of $\mathcal{E}(t)$, and (b) extremal values of $\mathcal{E}(\infty)$ as a function of Δ_z when $\omega_0 = 1.4\omega_{11}$. The cases of $\omega_0 = 1.2\omega_{11}$ (blue blocks) and $1.0\omega_{11}$ (green triangles) are also given in (b). The other parameters are the same as Fig. 2.

These diverse behaviors can be explained by the energy-spectrum feature. We see that two branches of bound states in the band gap region separate the energy spectrum into three regimes. When $\omega_0 \geq 2.8\omega_{11}$, no bound state is formed and $\mathcal{E}(t)$ damps to zero exclusively. When $1.1\omega_{11} \leq \omega_0 < 2.8\omega_{11}$, one bound state is formed and $\mathcal{E}(t)$ tends to a finite value. As long as two bound states are present when $\omega_0 < 1.1\omega_{11}$, $\mathcal{E}(t)$ exhibits a lossless Rabi-like oscillation. The matching of the long-time behaviors in the three regimes with the analytical result in Eq. (8) verifies the distinguished role played by the bound states and non-Markovian effect in the charging performance of our scheme. The result reveals that we can properly choose the working frequency ω_0 of the QB and charger such that a remote charging to the QB is dynamically realized without resorting to the direct charger-QB interactions.

To verify the tolerance of the scheme to the increase of the charger-QB distance, we plot in Fig. 3 the eigenenergies of bound states, the long-time evolution of $\mathcal{E}(t)$, and the extremal values of $\mathcal{E}(\infty)$ as a function of Δ_z. The result for $\omega_0 = 1.4\omega_{11}$ shows that when $\Delta_z < 0.3\lambda_{11}$, one bound state is present and $\mathcal{E}(t)$ evolves to a finite Δ_z-dependent value. When $0.3\lambda_{11} \leq \Delta_z < 1.7\lambda_{11}$, two bound states are formed and $\mathcal{E}(t)$ tends to a persistent Rabi-like oscillation in a frequency $E^{\rm b}_{+} - E^{\rm b}_{-}$ characterizing a coherent energy exchange between the QB and the charger. When $\Delta_z \geq 1.7\lambda_{11}$, the eigenenergies of the two bound states become degenerate and $\mathcal{E}(t)$ tends to zero according to Eq. (8). Therefore, our charging scheme can work well when Δ_z is freely changed within the range from $0.3\lambda_{11}$ to $1.7\lambda_{11}$. The result in Fig. 3(b) also reveals that the maximal stored energy $\max \mathcal{E}(\infty)$ can be enhanced by properly choosing the working frequency ω_0 of the QB and the charger. This gives us sufficient room to optimize the performance of our remote-charging scheme. Note that the presence of two bound states in Fig. 4(a) and nonzero $\mathcal{E}(\infty)$ are not always accompanied by an extractable energy of the QB described by the nonzero ergotropy $\mathcal{W}(\infty)$, see Fig. 4(b). Via optimizing ω_0, we plot $\max_{\omega_0} \mathcal{W}(\infty)$ in Fig. 4(c) in different Δ_z, which exhibits a good performance of our QB in quite a large charger-QB distance regime.

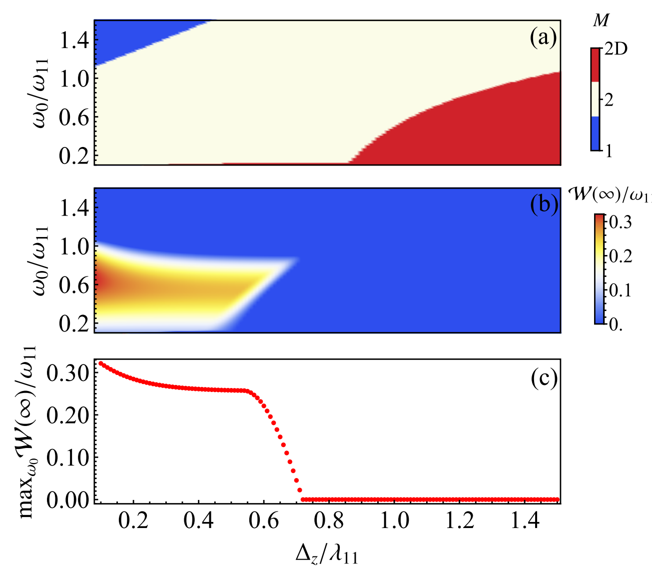

FIG. 4. (a) Number M of the bound states, (b) maxima of the ergotropy $\mathcal{W}(\infty)$ in different ω_0, and (c) maxima of the ergotropy $\max_{\omega_0} \mathcal{W}(\infty)$ via optimizing ω_0 as a function of Δ_z. 2D means two degenerate bound states. The other parameters are the same as Fig. 2.

Discussions and conclusion.—Although only the rectangular hollow metal waveguide is considered, our result is applicable to other systems. Applying it in the surface plasmon polariton waveguide [88,89], we are even not bothered to put the charger and the QB inside the waveguide, which reduces the experimental difficulty. Current experimental advances of the waveguide QED provide support for our QB scheme [90–93]. Especially, based on the solid-state emitter, i.e., nitrogen-vacancy centers (NVC), a series of integrated platforms have ensured the feasibility of our proposal, such as microwave coplanar [94], silicon nitride [95], laser-written [96], and nanowire [97] waveguides. The long-range interaction of NVCs in magnon waveguide QED has been proposed [98]. In our setup, proper values of the working frequency ω_0 and the charger-QB distance Δ_z are required. NVC is of the capability for coherent manipulation by either light or microwave. In the optical frequency regime, we can focus on the working frequency $\lambda_0 = 637$ nm. To match $\omega_0 = 1.0\omega_{11}$ for the maximal stored energy $\mathcal{E} = 0.6\omega_{11}$, the transverse lengths of waveguide are set as $a = b = 0.45$ μm. To match $\omega_0 = 0.6\omega_{11}$ for the maximal extractable energy $\mathcal{W} = 0.3\omega_{11}$, they are set as $a = b = 0.27$ μm. In the microwave frequency regime, the working frequency of NVC with zero-field splitting 2.87 GHz has a tremendous tunability by an external field. Generally speaking, the effective charging distance could range from 1 μm in the optical waveguide to 0.1 m in the microwave one, which are far beyond the typical length of the dipole-dipole coupling. The bound state and its distinguished role in the non-Markovian dynamics have been observed in both photonic crystal [99] and ultracold-atom [100,101] systems, which sets a solid foundation to the realizability of our scheme. Also, presenting a scalable and tunable framework, our scheme is easily extended to the multi-TLS scenarios.

In conclusion, we have proposed a remote-charging scheme via coupling the QB and the charger to a rectangular hollow metal waveguide. It is found that, contrary to one's belief that decoherence generally leads to the aging of a QB, the non-Markovian decoherence induced by the EMF in the waveguide can realize a persistent Rabi-like energy exchange between the QB and the charger. Our analysis reveals that such an ideal charging performance is caused by the formation of two bound states in the energy spectrum of the total system consisting of the QB, the charger, and the EMF in the waveguide. Without resorting to direct charger-QB coupling, our scheme avoids the insufficient-charging difficulty encountered in the case when the charger-QB coupling is weak. It provides a useful path to the practical realization of antiaging QB.

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*binzhou@hubu.edu.cn
†ywl@wipm.ac.cn
‡anjhong@lzu.edu.cn

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