

Inhomogeneous mass trap for dark-state polaritons in atomic media

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The generation of a trapping potential for dark-state polaritons in a two-dimensional electromagnetically induced transparency system is theoretically studied. We show that such a trap can arise from a spatially inhomogeneous effective mass of the dark-state polariton. Because this mass inhomogeneity can be engineered by tuning the parameters of the control fields, the motion, spatial profile, and coherent behavior of bound dark-state spatial modes can be tailored accordingly. Our results enable spatial controls of optical information and provide a possible route toward realizing Bose-Einstein condensation of dark-state polaritons in a trapping potential.

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Electromagnetically induced transparency (EIT) enables coherent control of light-matter interactions [1], with striking effects including slow light [2] and optical quantum memory [3–8]. Recent advances in structured light have expanded EIT functionalities, such as optical vortex exchange [9–11]. In counterpropagating schemes, stationary light pulses (SLPs) and dark-state polaritons (DSPs) with vanishing group velocity are generated [12–17]. DSPs have controllable effective mass and kinetics [18–21], enabling long interaction times for quantum information processing and Rydberg interactions [19,22–24]. However, most studies of SLPs assume uniform control fields [13,14,18–21,25,26], leaving structured-control-field effects largely unexplored [27–31]. Here, we show that structured control fields generate scalar and vector potentials for DSPs, producing coherent confinement and spatially varying attenuation that lead to bound spatial modes. Unlike methods based on elaborate optical trajectories [32] or rotations [20], our purely optical approach is more flexible. Our results pave the way for Bose-Einstein condensation of DSPs in trapping potentials [19,24,33,34]. We study a two-dimensional counterpropagating three-level Λ -type EIT [1,3,4,12,35,36], as shown in Figs. 1(a) and 1(b). The control fields (yellow arrows) couple the $|2\rangle \rightarrow |3\rangle$ transition with Rabi frequency $\Omega_c^{F(B)}$ and detuning $\Delta_c^{F(B)}$. The probe fields (blue arrows) drive the $|1\rangle \rightarrow |3\rangle$ transition with $\Omega_p^{F(B)}$ and $\Delta_p^{F(B)} = \Delta_c^{F(B)}$. Here, F and B denote forward- and backward-propagating fields along the y axis. The governing optical Bloch equations (OBEs) are derived under the condition

$|\Omega_c^{F(B)}| \gg |\Omega_p^{F(B)}|$. This results in the perturbative regime $\rho_{11} \approx 1 \gg |\rho_{a1}| \gg |\rho_{ab}|$ for $a, b > 1$, where ρ_{ab} represents the atomic coherence between states $|a\rangle$ and $|b\rangle$. The OBEs of the perturbative coherences read [12–14,20,21,30,31,37,38]

$$\partial_t \rho_{21} = \frac{i}{2} (\Omega_c^{F*} \rho_{31}^F + \Omega_c^{B*} \rho_{31}^B), \quad (1)$$

$$\partial_t \rho_{31}^F = \frac{i}{2} \Omega_p^F + \frac{i}{2} \Omega_c^F \rho_{21} - \left(\frac{\Gamma}{2} + i\Delta_p^F \right) \rho_{31}^F, \quad (2)$$

$$\partial_t \rho_{31}^B = \frac{i}{2} \Omega_p^B + \frac{i}{2} \Omega_c^B \rho_{21} - \left(\frac{\Gamma}{2} + i\Delta_p^B \right) \rho_{31}^B, \quad (3)$$

$$\frac{1}{c} \partial_t \Omega_p^F + \partial_y \Omega_p^F = i\eta \rho_{31}^F + \frac{i}{2k_p} \partial_z^2 \Omega_p^F, \quad (4)$$

$$\frac{1}{c} \partial_t \Omega_p^B - \partial_y \Omega_p^B = i\eta \rho_{31}^B + \frac{i}{2k_p} \partial_z^2 \Omega_p^B. \quad (5)$$

Here, ρ_{21} is the ground-state coherence between $|2\rangle$ and $|1\rangle$. $\rho_{31}^{F(B)}$ is the coherence between $|3\rangle$ and $|1\rangle$ driven by the forward (backward) probe field $\Omega_p^{F(B)}$. Γ is the spontaneous decay rate of the state $|3\rangle$. $\eta = \frac{\Gamma \xi}{2L}$ is the light-matter coupling constant, where ξ is the optical depth and L is the medium length along the y axis. $k_p = 2\pi/\lambda_p$ is the wave number of probe fields, λ_p is the probe wavelength, and c is the speed of light in vacuum. For simplicity, we set $\Delta_{p(c)}^F = \Delta_{p(c)}^B = \Delta_{p(c)}$ and two-photon resonance in the following analysis.

By neglecting time derivatives in Eqs. (2)–(5), $\rho_{31}^{F(B)}$ are expressed in terms of $\Omega_p^{F(B)}$. Substituting these back iteratively and retaining second-order spatial derivative terms, the truncated $\rho_{31}^{F(B)}$ are then inserted into Eq. (1), which yields a Schrödinger-like equation of ρ_{21} carrying effective charge one [30,31,37]:

$$i\hbar \partial_t \rho_{21} = \frac{(\frac{\hbar}{i} \partial_z - A_z)^2}{2M_z} \rho_{21} + \frac{(\frac{\hbar}{i} \partial_y - A_y)^2}{2M_y} \rho_{21} + U \rho_{21}, \quad (6)$$

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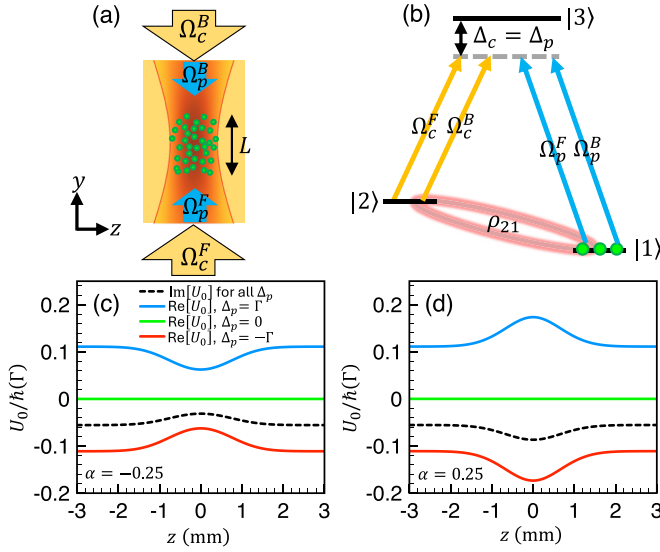


FIG. 1. (a) Two-dimensional counterpropagating EIT system. Green dots represent the atomic medium of length L . Yellow arrows denote the counterpropagating biased Gaussian control fields $\Omega_c^{F(B)}$, and blue arrows indicate the probe fields $\Omega_p^{F(B)}$. (b) Three-level Λ -type scheme, where $\Omega_p^{F(B)}$ and $\Omega_c^{F(B)}$ drive the transitions $|1\rangle \rightarrow |3\rangle$ and $|2\rangle \rightarrow |3\rangle$ with detuning $\Delta_p = \Delta_c$. Inhomogeneous mass trap potential $U_0/\hbar$ with (c) $\alpha = -0.25$ and (d) $\alpha = 0.25$. Red solid, green solid, and blue solid lines depict the real part of $U_0/\hbar$ for $\Delta_p = -\Gamma$, 0 , and Γ , respectively. Black dashed line illustrates the imaginary part of $U_0/\hbar$. Other parameters are $(w_0, \phi, \xi) = (1 \text{ mm}, 0, 80)$.

where $\hbar$ is the reduced Planck constant. The effective mass $M_{z(y)}$, the $z(y)$ component of the synthetic vector potential $A_{z(y)}$, and the synthetic scalar potential U for the ground-state coherence ρ_{21} are

$$M_z = \frac{\hbar k_p}{V_F + V_B}, \quad (7)$$

$$M_y = \frac{\hbar \eta}{2(2\Delta_p - i\Gamma)(V_F + V_B)}, \quad (8)$$

$$A_z \approx \frac{iM_z}{2\eta k_p} (\Omega_c^{F*} \partial_z \Omega_c^F + \Omega_c^{B*} \partial_z \Omega_c^B), \quad (9)$$

$$A_y \approx M_y (V_B - V_F), \quad (10)$$

$$U \approx \frac{i\hbar}{2\eta} (\Omega_c^{B*} \partial_y \Omega_c^B - \Omega_c^{F*} \partial_y \Omega_c^F) - \frac{A_z^2}{2M_z} - \frac{A_y^2}{2M_y} + \frac{\hbar \partial_z A_z}{2iM_z} + \frac{\hbar \partial_y A_y}{2iM_y}, \quad (11)$$

where $V_{F(B)} = \frac{|\Omega_c^{F(B)}|^2}{2\eta}$ denote the forward (backward) EIT group velocity. We aim to generate a synthetic trapping potential for ρ_{21} by employing biased Gaussian beams as the control fields:

$$\Omega_c^F(\vec{r}) = \Omega \{1 + \alpha G(\vec{r}) e^{-i\Phi(\vec{r})}\}, \quad (12)$$

$$\Omega_c^B(\vec{r}) = \Omega \{1 + \alpha G(\vec{r}) e^{i[\Phi(\vec{r}) + 2\phi]}\}. \quad (13)$$

Here, each Gaussian beams is biased by a uniform field Ω . The parameter α represents the fractional contribution

of the Gaussian beam component with the field amplitude $G(\vec{r}) = \frac{1}{\sqrt{1+y^2/L_r^2}} \exp[-\frac{z^2}{w_0^2(1+y^2/L_r^2)}]$ and the phase $\Phi(\vec{r}) = \frac{z^2}{w_0^2(1+y^2/L_r^2)} \frac{y}{L_r}$. $L_r = \pi w_0^2/\lambda_c$ is the Rayleigh range, w_0 is the beam waist radius, and λ_c is the wavelength of control fields. 2ϕ denotes the phase shift between the backward Gaussian beam and the backward uniform field. In the following, we perform numerical calculations using the common parameters: $(\alpha, \Gamma, \Omega, \lambda_c, \lambda_p, L) = (-0.25, 2\pi \times 6 \text{ MHz}, 3\Gamma, 780 \text{ nm}, 780 \text{ nm}, 5 \text{ mm})$.

We consider the situation $L_r \gg L$ where Eqs. (7)–(13) are uniform along y inside the Rayleigh range, and Eq. (6) can be approximated as

$$\begin{aligned} i\hbar \partial_t \rho_{21} &\approx -\frac{\hbar^2}{2M_z} \partial_z^2 \rho_{21} - \frac{\hbar^2}{2M_y} \partial_y^2 \rho_{21} + \frac{i\hbar A_y}{M_y} \partial_y \rho_{21}, \\ M_z &= \frac{\hbar k_p \eta}{\Omega^2 \left[1 + 2\alpha \cos^2(\phi) e^{-\frac{z^2}{w_0^2}} + \alpha^2 e^{-2\frac{z^2}{w_0^2}} \right]}, \\ M_y &= \frac{\eta M_z}{2(2\Delta_p - i\Gamma)k_p}, \\ A_y &= \frac{-2\alpha M_y \Omega^2 \sin^2(\phi)}{\eta} e^{-\frac{z^2}{w_0^2}}. \end{aligned} \quad (14)$$

Since $\Omega_c^{F(B)}$ depends only on z within L_r and the SLP is confined by the boundary conditions $\rho_{21} = 0$ at $y = \pm L/2$, the y -dependent solution is expected to take the form $\cos(k_y y)$. We therefore apply the standard gauge transformation for A_y and substitute $\rho_{21} = \psi_{nm}(z) \cos[(m + \frac{1}{2})k_y y] e^{\frac{i}{\hbar} A_y y - i\nu_{nm} t}$ and $k_y = 2\pi/L$ into Eq. (14) and arrive at $-\frac{\hbar^2}{2M_z} \partial_z^2 \psi_{nm} + U_m \psi_{nm} = \hbar \nu_{nm} \psi_{nm}$, where the effective potential reads

$$U_m \approx \left(m + \frac{1}{2}\right)^2 \frac{\hbar^2 k_y^2}{2M_y} - \frac{A_y^2}{2M_y}. \quad (15)$$

Since U_m results from the inhomogeneity of M_y , we term it as the inhomogeneous mass trap (IMT) for the coherence ρ_{21} . Figure 1(c) illustrates the IMT potential $U_0/\hbar$ for a negative parameter $\alpha = -0.25$, and Fig. 1(d) shows IMT for a positive $\alpha = 0.25$. In both panels, the real part of IMT $\text{Re}[U_0/\hbar]$ is shown for three different probe detunings: $\Delta_p = -\Gamma$ (red solid line), $\Delta_p = 0$ (green solid line), and $\Delta_p = \Gamma$ (blue solid line). The remaining parameters are set to $(w_0, \phi, \xi) = (1 \text{ mm}, 0, 80)$ and the longitudinal quantum number $m = 0$. Given that a convex potential, e.g., the blue solid line in Fig. 1(c) and the red solid line in Fig. 1(d), typically generates bound states, one might expect the bound-state condition to be $\alpha \Delta_p < 0$. However, our numerical solutions of the OBE show that centrally confined DSP profiles occur only when $\alpha < 0$. We find that this apparent discrepancy can be explained by the non-Hermiticity of the IMT. The imaginary part of the IMT $\text{Im}[U_0]$ (black dashed line for all Δ_p) behaves like an incoherent spatial filter that a negative α gives a concave $\text{Im}[U_0]$ [Fig. 1(c)] and leads to stronger attenuation of ρ_{21} outside the Gaussian beam waist than inside, whereas a positive α causes the opposite behavior [Fig. 1(d)]. By calculating the trap depth

$D_m = U_m|_{z=\infty} - U_m|_{z=0}$, we obtain

$$D_m = -\frac{\hbar\alpha[\alpha + 2\cos^2(\phi)](2\Delta_p - i\Gamma)(1 + 2m)^2 k_y^2 \Omega^2}{4\eta^2} + \frac{\hbar(2\Delta_p + i\Gamma)\alpha^2 \Omega^2 \sin^4(\phi)}{(4\Delta_p^2 + \Gamma^2)[1 + 2\alpha\cos^2(\phi) + \alpha^2]}. \quad (16)$$

Equation (16) confirms that $\alpha < 0$ makes $\text{Im}[D_m] < 0$ and incoherently results in a centrally confined profile. Consequently, the combination of $\alpha < 0$ and $\Delta_p > 0$ will coherently generate a SLP bound state as indicated by the blue line in Fig. 1(c). As $\Omega_c^{F(B)}$ increases radially along the z axis, SLPs near the center experience weaker attenuation and more dispersion than those at the edges. In cold-atom media, the use of $\Delta_p^F \neq \Delta_p^B$ suppresses the SLP leakage [14]. Equation (15) can then be generalized for $\phi = 0$ by replacing $2\Delta_p$ with $\Delta_p^F + \Delta_p^B$. In the following, we focus on the case $\alpha < 0$, which produces the desired potential well and supports bound DSP states. We use the Taylor's expansion of $U_m(z)$ at $z = 0$ and obtain the Schrödinger equation for the quantum harmonic oscillator with complex frequency [39]:

$$-\frac{\hbar^2}{2M_z^c} \partial_z^2 \psi_{nm} + \frac{M_z^c \omega_m^2}{2} z^2 \psi_{nm} = \left(n + \frac{1}{2}\right) \hbar \omega_m \psi_{nm}. \quad (17)$$

Its eigenvalues $(n + \frac{1}{2})\omega_m = \nu_{nm} - \delta_m$ and eigenfunctions are

$$\delta_m = \left(m + \frac{1}{2}\right)^2 \frac{\hbar k_y^2}{2M_y^c} - \frac{(A_y^c)^2}{2\hbar M_y^c}, \quad (18)$$

$$\omega_m^2 = \left(m + \frac{1}{2}\right)^2 \frac{2\hbar^2 k_y^2 (-\alpha)[\alpha + \cos^2(\phi)]}{M_y^c M_z^c w_0^2 [1 + 2\alpha\cos^2(\phi) + \alpha^2]} + \frac{2(A_y^c)^2 [1 + \alpha\cos^2(\phi)]}{M_y^c M_z^c w_0^2 [1 + 2\alpha\cos^2(\phi) + \alpha^2]}, \quad (19)$$

$$\psi_{nm} = c_{nm} H_n \left(\sqrt{\frac{M_z^c \omega_m}{\hbar}} z \right) \exp \left(-\frac{M_z^c \omega_m}{2\hbar} z^2 \right), \quad (20)$$

where $M_y^c = \frac{\hbar\eta^2(2\Delta_p + i\Gamma)}{2\Omega^2[1 + 2\alpha\cos^2(\phi) + \alpha^2](4\Delta_p^2 + \Gamma^2)} = |M_y^c|e^{i\theta}$ and $M_z^c = \frac{\hbar k_p \eta}{\Omega^2[1 + 2\alpha\cos^2(\phi) + \alpha^2]}$ are the central effective masses, and $A_y^c = \frac{-\hbar\alpha\eta(2\Delta_p + i\Gamma)\sin^2(\phi)}{[1 + 2\alpha\cos^2(\phi) + \alpha^2](4\Delta_p^2 + \Gamma^2)}$ is the central vector potential at $z = 0$. $|M_y^c| = \frac{\hbar\eta^2}{2\Omega^2[1 + 2\alpha\cos^2(\phi) + \alpha^2]\sqrt{4\Delta_p^2 + \Gamma^2}}$ is the mass module, $\theta = \cos^{-1}(\frac{2\Delta_p}{\sqrt{4\Delta_p^2 + \Gamma^2}})$ is the mass angle, and c_{nm} is the normalization constant. Equations (14)–(16) suggest the physical picture that two counterpropagating biased Gaussian fields $\Omega_c^{F(B)}$, on one hand, constitute a cavity along the y direction and, on the other hand, generate a harmonic trap along the z direction. Figure 2 compares the analytic solution of Eq. (20) (blue dotted line) with results from two numerical methods [30,31]: steady-state solution of OBE, i.e., Eqs. (1)–(5) (green solid line) and the eigenfunction of Eq. (6) (red dash-dotted line). As higher-energy states attenuate faster than the ground state, the steady-state solution of the OBE is expected to converge to the ground state. Figures 2(a) and 2(b) illustrate the probability densities $|\psi_{00}|^2$ and $|\psi_{10}|^2$, respectively. These probability density plots are taken as cross sections along

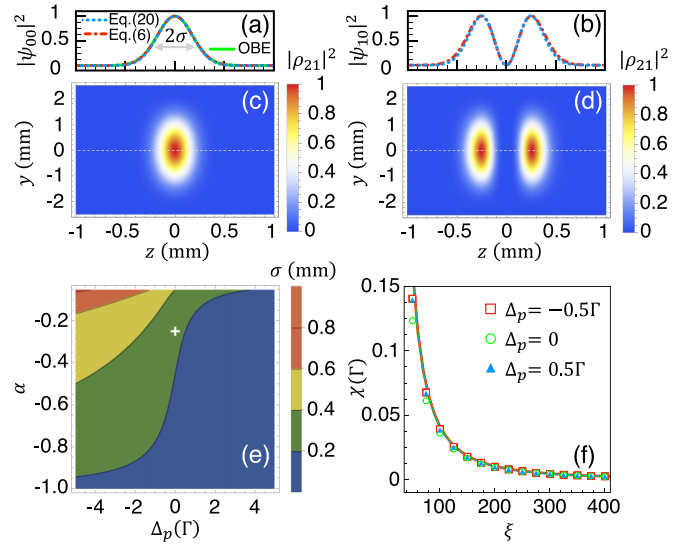


FIG. 2. (a) $|\psi_{00}|^2$ and (b) $|\psi_{10}|^2$ are the cross sections along white dashed lines in (c) ground state and (d) first excited state $|\rho_{21}|^2$, respectively. Green solid line is the steady-state $|\rho_{21s}|^2$ from OBE, red dash-dotted lines are the eigenfunction of Eq. (6), and blue dashed lines represent Eq. (20). (e) Half width σ of $|\psi_{00}|^2$. The white cross pinpoints parameters $(\Delta_p, w_0, \phi, \xi) = (0, 1 \text{ mm}, 0, 80)$ used in panels (a)–(d). (f) Red squares, green circles, and blue triangles are χ from the numerical solution of OBE with above parameters except $\Delta_p = -0.5\Gamma, 0$, and 0.5Γ , respectively. Color solid lines represent Eq. (22).

the white dashed lines shown in Fig. 2(c) (ground state) and Fig. 2(d) (first excited state) from the numerical diagonalization of Eq. (6). Both figures confirm a good agreement between the analytic and numerical solutions with parameters $(\Delta_p, w_0, \phi, \xi) = (0, 1 \text{ mm}, 0, 80)$.

To show the confinement and the non-Hermiticity, we derive $|\psi_{00}|^2 \propto \exp(-z^2/\sigma^2 - \chi t)$ for $\phi = 0$, where the half width σ and the decay rate $\chi = 2\text{Im}[\nu_{00}]$, i.e., twice the imaginary part of ν_{00} , are

$$\sigma = \sqrt{\frac{w_0}{k_y}} \left[\frac{2\eta(1+\alpha)}{-\alpha k_p (2\Delta_p + \sqrt{\Gamma^2 + 4\Delta_p^2})} \right]^{\frac{1}{4}}, \quad (21)$$

$$\chi = \frac{\Gamma k_y^2 \Omega^2 (1+\alpha)^2}{2\eta^2} + \frac{k_y \Gamma \Omega^2 \sqrt{-\alpha(1+\alpha)^3}}{w_0 \sqrt{2k_p \eta^3 (2\Delta_p + \sqrt{\Gamma^2 + 4\Delta_p^2})}}. \quad (22)$$

Based on Fig. 1(c), $\Delta_p > 0$ generates a deeper IMT than $\Delta_p \leq 0$ does. This deeper well confines ρ_{21} more tightly, as illustrated in Fig. 2(e). $\sigma \propto \eta^{1/4} \propto (\xi/L)^{1/4}$ shows a weak dependence on ξ/L . For example, σ increases by less than a factor of 2 as ξ changes from 80 to 1000 for a given L . The decay rate χ reflects the non-Hermitian nature of the system. Since the first term in Eq. (22) dominates and the dependence on Δ_p is negligible, $\chi \propto (\xi/L)^{-2}$ indicates that a larger ξ/L suppresses loss, making the system more Hermitian-like. Figure 2(f) shows this trend and the excellent agreement between Eq. (22) (solid lines) and numerical results for $\Delta_p = -0.5\Gamma$ (red squares), 0 (green circles), and 0.5Γ (blue triangles), using the same parameters as in Figs. 2(a)–2(d).

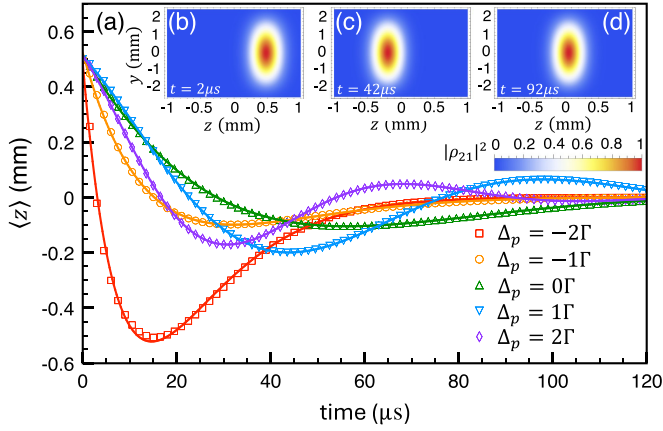


FIG. 3. (a) Coherent state oscillation of the expectation value $\langle z \rangle$ for $\Delta_p = -2\Gamma$ (red squares), $\Delta_p = -1\Gamma$ (orange circles), $\Delta_p = 0\Gamma$ (green triangles), $\Delta_p = 1\Gamma$ (blue inverted triangles), and $\Delta_p = 2\Gamma$ (purple diamonds). The insets depict the snapshot of $|\rho_{21}(z, y)|^2$ for $(\Delta_p, w_0, \phi, \xi) = (1\Gamma, 1.5 \text{ mm}, 0, 80)$ at (b) $t = 2 \mu\text{s}$, (c) $t = 42 \mu\text{s}$, and (d) $t = 92 \mu\text{s}$.

We next examine coherent effects by analyzing ρ_{21} in a coherent state and demonstrate that the IMT goes beyond a simple spatial filter. We numerically solve OBE with the parameters $(w_0, \phi, \xi) = (1.5 \text{ mm}, 0, 80)$ to investigate the dynamics of the coherent-state $\rho_{21} = c_{00} \exp[-\frac{M_z^2 \omega_m}{2\hbar}(z - z_0)^2]$. Here, z_0 represents the peak position that is intentionally shifted from the IMT center. This initial wave packet is prepared via the EIT light storage [3–7]. We calculate the subsequent expectation value of its position $\langle z(t) \rangle = [\int z |\rho_{21}(\vec{r}, t)|^2 d^2\vec{r}] / [\int |\rho_{21}(\vec{r}, t)|^2 d^2\vec{r}]$. Figure 3(a) presents the evolution of $\langle z(t) \rangle$ for $\Delta_p = -2\Gamma$ (red squares), -1Γ (orange circles), 0 (green triangles), 1Γ (blue-inverted triangles), and 2Γ (purple diamonds). Typical damped oscillations are observed. For $\Delta_p = 1\Gamma$, snapshots of $|\rho_{21}(\vec{r}, t)|^2$ at $t = 2, 42$, and $92 \mu\text{s}$ are shown in Figs. 3(b)–3(d), revealing back-and-forth oscillation of the wave packet. This behavior is quantified by fitting the numerical results to $e^{-\kappa t} \cos(ft)$. Each solid line in Fig. 3(a) shows the optimal fitting curve with the fitting parameters $(f, \kappa) = (21.1 \text{ kHz}, 84.7 \text{ kHz})$ for $\Delta_p = -2\Gamma$, $(f, \kappa) = (27.4 \text{ kHz}, 54.9 \text{ kHz})$ for $\Delta_p = -1\Gamma$, $(f, \kappa) = (29 \text{ kHz}, 28 \text{ kHz})$ for $\Delta_p = 0$, $(f, \kappa) = (58.1 \text{ kHz}, 20.8 \text{ kHz})$ for $\Delta_p = 1\Gamma$, and $(f, \kappa) = (84.1 \text{ kHz}, 33.7 \text{ kHz})$ for $\Delta_p = 2\Gamma$. The fitted oscillation frequency f shows a good agreement with the theoretical prediction of the real part of the trap frequency [Eq. (19)] for $\Delta_p > 0$ giving $\text{Re}[\omega_0] = 66 \text{ kHz}$ for $\Delta_p = 1\Gamma$ and $\text{Re}[\omega_0] = 91.6 \text{ kHz}$ for $\Delta_p = 2\Gamma$. The under-damped oscillation occurs only for $\Delta_p > 0$, and the oscillation frequency increases with increasing Δ_p . Moreover, for $\Delta_p \leq 0$, the wave packet crosses $z = 0$ only once before returning to the origin. Both behaviors confirm the picture in Fig. 1(c); namely, the combination of $\alpha < 0$ and $\Delta_p > 0$ produces a deeper IMT than the case with $\Delta_p < 0$. Figures 2 and 3 show excellent agreement between numerical and theoretical results, as well as the remarkable controllability of attenuation and coherent dynamics in the non-Hermitian IMT.

We now study the effect of the Gaussian beam phase shift ϕ and calculate the ground-state position expectation

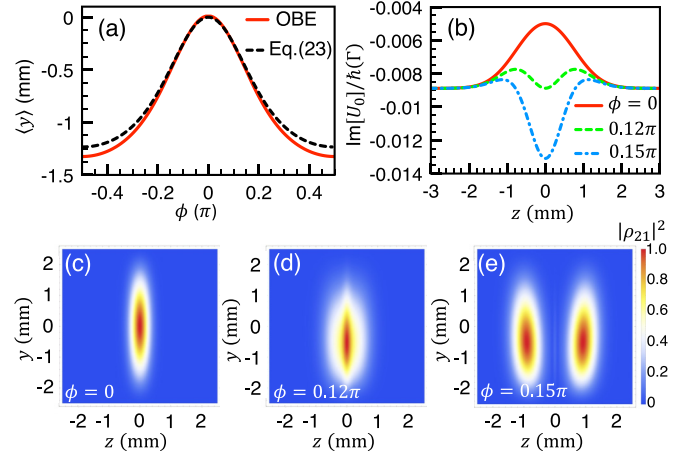


FIG. 4. (a) ϕ -dependent expectation value $\langle y \rangle$ at $z = 0$. The red solid line shows the numerical solution of the OBE, and the black-dashed line represents Eq. (23). (b) $\text{Im}[U_0]/\hbar$ for $\phi = 0$ (red solid line), $\phi = \phi_c = 0.12\pi$ (green dashed line), and $\phi = 0.15\pi$ (blue dash-dotted line). (c) $|\rho_{21}(z, y)|^2$ for $\phi = 0$, (d) $\phi = 0.12\pi$, and (e) $\phi = 0.15\pi$. Other parameters are $(\Delta_p, w_0, \xi) = (1\Gamma, 1.5 \text{ mm}, 200)$.

value $\langle y \rangle = (\int y |\rho_{21}|^2 dy) / \int |\rho_{21}|^2 dy$, where the ground-state $|\rho_{21}|^2 \propto \exp\{-\frac{2}{\hbar} \text{Im}[A_y]y\}$. The exponent reveals a key feature that A_y can push or pull DSPs along the y direction. For instance, when $\frac{2}{\hbar} \text{Im}[A_y] > 0$, the vector potential introduces an asymmetry and pushes the peak of the wave function toward the region $y < 0$. By evaluating $\langle y \rangle$, we obtain the analytical expression for the ϕ -dependent displacement:

$$\langle y \rangle = \frac{3k_1^2 + k_y^2}{k_1(k_1^2 + k_y^2)} - \frac{\pi}{k_y} \coth\left(\pi \frac{k_1}{k_y}\right), \quad (23)$$

where $k_1 = \frac{2}{\hbar} \text{Im}[A_y]$. Figure 4(a) presents the comparison between Eq. (23) (black dashed line) and the numerical $\langle y \rangle$ from OBE (red solid line) with the parameters $(\Delta_p, w_0, \xi) = (1\Gamma, 1.5 \text{ mm}, 200)$ at $z = 0$. The good agreement between the analytical result and the numerical solution confirms the predicted ϕ -induced negative displacement of the ground-state wave packet. Figures 4(c), 4(d), and 4(e) display the spatial distribution of the steady-state $|\rho_{21}|^2$ for $\phi = 0, 0.12\pi$, and 0.15π , respectively. We note that as ϕ is increased, the wave packet not only undergoes a backward displacement but also exhibits a distinct splitting when $\phi > 0.12\pi$. The ϕ -induced splitting of ρ_{21} can be attributed to the nonuniform attenuation. Figure 4(b) illustrates the imaginary part of IMT for $\phi = 0$ (red solid line), $\phi = 0.12\pi$ (green dashed line), and $\phi = 0.15\pi$ (blue dash-dotted line). For $\phi = 0$, the concave profile of $\text{Im}[U_0]$ leads to weaker attenuation of ρ_{21} inside the Gaussian beam waist than outside. As ϕ increases, $\text{Im}[U_0]$ becomes convex, causing greater attenuation inside the beam waist than outside. This spatially varying attenuation results in the observed splitting of ρ_{21} . We determine the critical phase ϕ_c at which the splitting emerges by solving $\text{Im}[D_0] = 0$,

resulting in a well-defined threshold condition

$$\phi_c = \tan^{-1} \sqrt{\frac{-(2+\alpha)(1+\alpha)^2}{1+\alpha(1+\alpha)^2 - \sqrt{1 - \frac{4\alpha(2+\alpha)(1+\alpha)^2\eta^2}{k_y^2(\Gamma^2+4\Delta_p^2)}}}}. \quad (24)$$

Equation (24) yields the exact numerical value $\phi_c = 0.12\pi$ for the parameters used in Fig. 4, in agreement with the observed results.

In conclusion, we have shown a method to generate IMT for the ground-state coherence ρ_{21} using two counter-propagating biased Gaussian beams under the guidance of Schrödinger-like equation (6). This IMT scheme enables tun-

able control of SLP properties, including attenuation, spatial profile, and coherent dynamics, via parameters such as laser detuning, coupling field strength, Gaussian phase, and optical depth. These results confirm the EIT system as an effective platform for DSP manipulation.

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Data availability. There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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