Electromagnetically Induced Transparency in Circuit Quantum Electrodynamics with Nested Polariton States

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Quantum networks will enable extraordinary capabilities for communicating and processing quantum information. These networks require a reliable means of storage, retrieval, and manipulation of quantum states at the network nodes. A node receives one or more coherent inputs and sends a conditional output to the next cascaded node in the network through a quantum channel. Here, we demonstrate this basic functionality by using the quantum interference mechanism of electromagnetically induced transparency in a transmon qubit coupled to a superconducting resonator. First, we apply a microwave bias, i.e., drive, to the qubit-cavity system to prepare a \(\Lambda\)-type three-level system of polariton states. Second, we input two interchangeable microwave signals, i.e., a probe tone and a control tone, and observe that transmission of the probe tone is conditional upon the presence of the control tone that switches the state of the device with up to 99.73% transmission extinction. Importantly, our electromagnetically induced transparency scheme uses all dipole allowed transitions. We infer high dark state preparation fidelities of \(> 99.39\%\) and negative group velocities of up to \(-0.52 \pm 0.09\, \text{km/s}\) based on our data.

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Controllable interaction between electromagnetic quanta and discrete levels in a quantum system, i.e., light matter interaction, is the key to quantum information storage and processing in a quantum network [1,2]. Consider a three-level atomic system driven by two coherent electromagnetic waves. The destructive interference between the two excitation pathways creates a transparency window for one of the drive fields and switches the system into a “dark state.” This phenomenon is called electromagnetically induced transparency (EIT) [3]. Recently, EIT has been harnessed for implementing different building blocks of a quantum network, such as all-optical switches and transistors [4–8], quantum storage devices [9–13], and conditional phase shifters [14–18]. Despite this remarkable success, utilizing EIT and related effects at the single-photon and single-atom level with highly scalable devices is a formidable challenge that prevents realization of a practical quantum network [19]. A promising solution is to extend these techniques to the microwave domain using superconducting quantum circuits that are both scalable and enable deterministic placement of long-lived artificial atoms for the network nodes [20–23].

To this end, three-level superconducting artificial atoms have been used to demonstrate coherent population trapping [24] and Autler-Townes splittings (ATS) [25–31]. However, conclusive evidence of EIT in these simple systems eluded researchers, as it is difficult to find a superconducting quantum system with metastable states and lifetimes that satisfy its stringent requirements [32–35]. Recently, progress has been made in a circuit quantum electrodynamics (QED) system that exploits qubit coupling to a single-mode cavity [36]. In that experiment, one leg of the \(\Lambda\)-type system is dipole forbidden, requiring that it be driven with a two-photon transition. The small photon scattering cross section of this two-photon transition hinders applications such as single-atom quantum memory [37], all-optical switching and routing of a single photon gated by another single photon [5], single photon-photon cross phase modulation [29], and vacuum-induced transparency [38]. On the other hand, high scattering cross sections have been observed in a dipole allowed transition of an artificial atom coupled to a one-dimensional waveguide [39]. Thus, implementing a \(\Lambda\)-type system with all dipole allowed transitions in a circuit QED system is highly desirable for building a quantum network with microwave photons.

In this Letter, we report the first observation of EIT using all dipole allowed transitions in a \(\Lambda\)-type system implemented with a superconducting quantum circuit. Our scheme is based on a theoretical proposal [40] that utilizes polariton states generated with a rf biased two-level system coupled to a resonator. Here, we realized the polariton states in a transmon-cavity system and achieved a metastable state with a long lifetime. Moreover, we were able to tune the polariton states to establish a \(\Lambda\)-type system that can be driven with control and probe fields through dipole...
allowed transitions. Note that due to the transmission geometry of our cavity where nominally the signal is transmitted on resonance, the observed experimental signal is actually electromagnetically induced absorption (EIA). However, our EIA and conventional EIT have identical underlying physics of quantum interference. Conventional EIT spectra can be observed if a hanger resonator geometry is used. We retain the nomenclature identical underlying physics of quantum interference.

Our experiment is performed on a device that consists of a concentric transmon capacitively coupled to a λ/2 microstrip resonator with a coupling strength g/2π = 74 MHz, as shown in Fig. 1(a). The transmon comprises a single Al/AlOx/Al Josephson junction shunted by a capacitor consisting of a superconducting island and a surrounding ring. The Josephson junction is fabricated with an overlap technique [41]. The transmon has a resonance frequency ωd/2π = 5.648 GHz between its lowest two levels and an anharmonicity α/2π = 262.5 MHz. The coherence times are measured to be T1 = 35 μs and T2 = 22.5 μs. The fundamental mode of the resonator is at ωr/2π = 6.485 GHz with an internal quality factor Qi = (1.2 ± 0.2) × 10⁶ and a loaded quality factor Q = 7900 dominated by the strong coupling to the microwave feedline at the output port.

The transmon-cavity system is well in the dispersive regime with a dispersive shift χ/2π = 1.54 MHz. The eigenlevels are described by the dispersive Jaynes-Cummings ladder as shown in Fig. 1(b). The resonance frequencies are ωq − χ for the |g, 0⟩ ↔ |e, 0⟩ transition and ωq − 3χ for the |g, 1⟩ ↔ |e, 1⟩ transition, where |g, n⟩ denotes the qubit ground (excited) state with n photons in the resonator. The tilde indicates that these levels are singly dressed states; i.e., they are transmon states slightly dressed with resonator photons.

The polariton states are generated by injecting a strong microwave drive field through the input coupler to doubly dress the Jaynes-Cummings states. In particular, if the drive frequency ωd is in the so-called nesting regime ωq − 3χ < ωd < ωq − χ, the resulting eigenstates |2⟩ and |3⟩ will be nested in between the eigenstates |1⟩ and |4⟩ [42–44].

We use the set of polariton states |1⟩, |2⟩, and |3⟩ to form a Λ-type system [Fig. 2(d)]. In the driven two-level-system model, these polariton states can be approximated as

\[ |1⟩ = −\sin \left( \frac{θ_0}{2} \right) |e, 0⟩ + \cos \left( \frac{θ_0}{2} \right) |g, 0⟩, \]

\[ |2⟩ = \cos \left( \frac{θ_0}{2} \right) |e, 0⟩ + \sin \left( \frac{θ_0}{2} \right) |g, 0⟩, \]

\[ |3⟩ = −\sin \left( \frac{θ_1}{2} \right) |g, 1⟩ + \cos \left( \frac{θ_1}{2} \right) |e, 1⟩, \]

\[ |4⟩ = \cos \left( \frac{θ_1}{2} \right) |g, 1⟩ + \sin \left( \frac{θ_1}{2} \right) |e, 1⟩, \]

where the mixing angles θ0 and θ1 are given by tan(θ0) = Ωd/[(ωq − χ) − ωd] and tan(θ1) = Ωd/[(ωq − χ) − (ωq − 3χ)] [40].

Equation (1) shows that the |1⟩ ↔ |3⟩ and |2⟩ ↔ |3⟩ transitions are mainly cavitylike transitions, while |1⟩ ↔ |2⟩ is a qubitlike transition. These properties can be revealed by calculating the decay rate γij of the |i⟩ → |j⟩ transition, which can be approximated as γ31 = γq sin²(θ0 + θ1)/2, γ32 = γq cos²(θ0 + θ1)/2, and γ21 = γq cos²(θ0)/2, where γq is the cavity decay rate, and γq is the qubit decay rate [40]. Thus, the decay rate of the
$|3\rangle \rightarrow |1\rangle$ transition ($\gamma_{31}$) can be tuned to be comparable to the decay rate of the $|3\rangle \rightarrow |2\rangle$ transition ($\gamma_{32}$), while extending the metastable state lifetime ($1/\gamma_{21}$) even beyond the qubit lifetime. These two effects are key to achieve EIT in our superconducting circuit system.

We measure the transition frequencies between the polariton states by performing two-tone spectroscopy with a polariton drive and a weak probe field. The drive frequency and the probe power are fixed at $\omega_d/2\pi = 5.6466$ GHz and $P_p = -163.15$ dB m, respectively, while scanning the drive strength and the probe frequency. The probe transmission defined as the ratio of the probe output complex amplitude to the input complex amplitude $S_{21} = V_{out}/V_{in} = |S_{21}| e^{i\phi}$ was measured by a vector network analyzer (VNA). Our definition of $S_{21}$ includes all round-trip amplification and attenuation, where $\phi$ has been corrected for electric delay. As shown in Fig. 2(b), there are four transmission peaks near the resonator frequency. The four peaks correspond to, from low to high frequencies, $\omega_{23}$, $\omega_{13}$, $\omega_{24}$, and $\omega_{14}$, respectively [Fig. 2(c)], where $\omega_{ij}$ denotes the energy difference between the polariton states $|i\rangle$ and $|j\rangle$. The spacing between the first and second (first and third) transmission peaks, which corresponds to the splitting between levels $|1\rangle$ and $|2\rangle$ ($|3\rangle$ and $|4\rangle$), widens as the drive strength increases. This is consistent with the expected ac Stark shift drawn as the black dashed curves in Fig. 2(b). Another crucial feature of the spectrum is that as the polariton drive strength increases, the height of the $\omega_{23}$ and $\omega_{14}$ peaks decreases, while the height of the $\omega_{13}$ and $\omega_{24}$ peaks increases. This behavior agrees with the change of the transition probabilities between the polariton states predicted in Ref. [40].

In this system, EIT is demonstrated by a suppression of transmission for a weak probe field on resonance with one leg of a $\Lambda$ system, while a control field addressing the other leg [Fig. 2(d)]. The $\Lambda$ system is established by a polariton drive field with frequency $\omega_d/2\pi = 5.6466$ GHz and strength $\Omega_d/2\pi = 1.46$ MHz. The resultant $\Lambda$ levels have $\gamma_{31}/2\pi = 0.35$ MHz and $\gamma_{32}/2\pi = 0.47$ MHz, which are much larger than $\gamma_{21}/2\pi = 2.74$ kHz. The control field frequency $\omega_c/2\pi = 6.4828$ GHz and the probe strength $\Omega_p/2\pi = 62$ kHz are fixed, while we scan the control field strength $\Omega_c$ and the probe frequency $\omega_p$. The probe transmission ($S_{21}$) measured by the VNA is shown in Figs. 3(a) and 3(b). With our parameters, the theoretical condition of EIT is given by $\Omega_c/2\pi < \gamma_{31}/2\pi = 0.82$ MHz [black dash-dotted line in Fig. 3(a)] [40]. Under this condition, we observe a transmission suppression window around $\omega_p = \omega_{13}$ with the largest suppression 25.66 dB [Figs. 3(c) and 3(d)], which means about 99.73% of the power of the original transmitted probe field is suppressed. However, as the control field strength exceeds the
EIT boundary, the transmission for \( \omega_p > \omega_{13} \) in Fig. 3(a) is becoming smaller and completely disappears above \( \Omega_c/2\pi = 2.8 \text{ MHz} \) instead of changing to an ATS line shape. This behavior is most likely due to excess cavity population above a single photon, due to the strong control field.

Quantum interferences in a driven \( \Lambda \) system create a dark state, which is transparent to the probe field. The fidelity of the dark state preparation is an important metric for an EIT-based quantum memory [19]. With our experimental parameters, we inferred the dark state fidelity defined as [45]

\[
F_s = \sqrt{\langle D | D \rangle},
\]

\[
= \sqrt{\frac{1}{2} \left[ \cos 2\Theta (\rho_{11} - \rho_{22} - \sin 2\Theta (\rho_{21} + \rho_{12}) + (1 - \rho_{11}) \right] },
\]

where the dark state \( |D\rangle = \cos \Theta |1\rangle - \sin \Theta |2\rangle \) and the mixing angle \( \Theta = \tan^{-1}(\Omega_p/\Omega_c) \). The density matrix \( \rho \) is calculated by numerically solving a Lindblad master equation of a driven system, including decay rates \( \gamma_{ij} \) [46]. At the EIT boundary (\( \Omega_c/2\pi = 0.82 \text{ MHz} \), \( \Omega_c/2\pi = 0 \text{ MHz} \)), the dark state fidelity is calculated to be 99.39%. Note that we switched the role of the probe and the control fields to simulate the fidelity when the dark state is essentially the polariton [2] state and the main infidelity source is its decay rate \( \gamma_{21} \).

To confirm that the suppression of transmission is due to EIT as opposed to ATS, Akaikes-information-criterion-(AIC) based testing was performed. The AIC-based testing calculates the weight of each fitting model based on the goodness of the fitting with the constraint that the sum of the weights is unity [32]. Originally, the AIC-based testing was proposed to fit the susceptibility \( \chi_s \) [32]. To use this criterion, we derive the relationship between the measured \( S_{21} \) and a generic susceptibility \( \chi_s \) as [47]

\[
\ln(S_{21}) = \ln(|S_{21}|) + i\phi = i\frac{\omega_p L}{c} \left( 1 + \frac{1}{2} \chi_s \right) - \alpha_0 + i\phi_0,
\]

where \( L \) is the effective distance the microwave travels through the chip, \( c \) is the speed of light, \( \alpha_0 \) is the attenuation of the cables, and \( \phi_0 \) is a frequency-independent initial phase offset. For EIT, the susceptibility takes the form of the difference between two Lorentzians [32]

\[
\chi^\text{EIT}_s = \frac{A_+}{(\omega_p - \omega_+) - i(\Gamma_+/2)} - \frac{A_-}{(\omega_p - \omega_-) - i(\Gamma_-/2)},
\]

and for ATS, it takes the form of the sum

\[
\chi^\text{ATS}_s = \frac{A_1}{(\omega_p - \omega_1) - i(\Gamma_1/2)} + \frac{A_2}{(\omega_p - \omega_2) - i(\Gamma_2/2)},
\]

where \( \omega_j, A_j, \) and \( \Gamma_j \) are the center frequency, magnitude, and width of the \( j \)th Lorentzian, respectively. In comparison to Ref. [32], the different negative signs in front of the \( i(\Gamma_j/2) \) terms in Eqs. (4) and (5) are due to the transmission geometry of the circuit. The model functions for EIT or ATS are then obtained by substituting either \( \chi^\text{EIT}_s \) or \( \chi^\text{ATS}_s \) for the \( \chi_s \) in Eq. (3).

We fit the probe transmission \( S_{21} \) data to both EIT and ATS models to extract the AIC-based testing weights to validate that the observations were from EIT [32]. For each model, \( \ln(|S_{21}|) \) and \( \phi \) were fit simultaneously to assure the Kramers-Kronig relations. The transmission data at \( \Omega_c/2\pi = 0.82 \text{ MHz} \) and their fits of both models are shown in Figs. 4(a) and 4(b). Qualitatively, at this control field strength, the data fit significantly better to the EIT model.
model than to the ATS model. Furthermore, the weights of the EIT and ATS models for different control field strengths are plotted in Fig. 4(c). For control field strength $\Omega_c/2\pi < 0.2$ MHz, both the EIT and ATS weights approach 0.5 due to the presence of noise and the relatively small size of transmission suppressions. In the range of $0.2$ MHz $< \Omega_c/2\pi < 2.8$ MHz, the EIT weights are substantially larger than the ATS weights, indicating strong EIT signatures. The maximum EIT weight happens around $\Omega_c/2\pi = 0.82$ MHz, which is in agreement with the theoretical EIT boundary. For control field strength $\Omega_c > 2.8$ MHz, the control field excites resonator photons and drives the system out of the nesting regime. Therefore, there is neither EIT nor ATS characters and results in equal weights of 0.5.

We also investigated the backward light phenomenon due to the giant dispersion of EIT [36]. We calculated the time $\tau_g$ for the probe field to traverse the device at different control field strengths by using $\tau_g = -d\phi/d\omega_p$, where $\phi$ is obtained from the fittings of the EIT model [Fig. 5(a)]. The group velocity of the probe can then be calculated by $v_g = l/\tau_g$, where $l = 10.3$ mm is the distance between the input and output coupler of the device. The largest inferred negative group velocity is $v_g = -0.52 \pm 0.09$ km/s, further pushing the boundaries of slow light, compared to that reported in Ref. [36].

In conclusion, polariton states in the nesting regime were generated with a transmon circuit QED system. The transmission spectra were measured and agree with theoretical predictions. We utilized three levels of nested polariton states to form a A-type transition. A robust EIT signature with all dipole allowed transitions was observed in a superconducting system for the first time. Our results constitute an important step toward a scalable quantum network with propagating microwave photons.

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