Transient coherence of an EIT media under strong phase modulation

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A strong phase modulated coupling field leads to an amplitude modulation of a probe field in an electromagnetically induced transparency process. Time vs. detuning plots for different modulation frequencies reveals a transition between an adiabatic regime where a series of smooth pulses are created with a global phase dependent upon the detuning, and a non-adiabatic regime where a strong transient oscillating response is added to these pulses. In the extreme non-adiabatic regime, where the modulation frequency is higher than the transient decay time, a coherent interference pattern is revealed. Adding a magnetic field lifts the hyperfine level degeneracy, resulting in a separation of the original pulse to three pulses. Every pulse has now a different phase dependent upon the magnetic field causing an interference effect between the different magnetic level transients. We explore the dynamics of the magnetic and non magnetic cases and show its resemblance to the Landau-Zener theory. We also show that combining the global phase of the pulses with the transient interference allows for a wide magnetic sensing range without loosing the sensitivity of a single EIT line.

Electromagnetically induced transparency (EIT) is a coherent process, where a strong coupling field creates a narrow transmission band in the probe spectrum in an otherwise fully absorptive medium \([1, 2]\). The narrow linewidth of EIT makes it suitable for applications in many fields such as extreme slow light, quantum storage devices, non linear optics and high sensitivity magnetic sensors. On the other hand, this narrow linewidth directly limits the bandwidth of data that can be processed. A signal which has a broader bandwidth than the EIT linewidth will be absorbed spectrally, or will not be delayed in the time domain \([3]\). In terms of magnetic sensing this means that although a very high sensitivity is possible using EIT, it is a problem to probe with this sensitivity a broadband field. A possible solution to this problem may arise from two directions. One idea is to use a multimode EIT system where each EIT line has still a narrow bandwidth but spreading the signal across many systems allows for a broader signal. Such systems were devised spatially \([4]\) and spectrally \([5, 6]\) for larger data capacity as well as for broadband magnetic sensing \([7, 8]\). A different approach is to use dynamic EIT where the transient response may have a much broader bandwidth than steady state EIT. The transient response of an EIT media to a sudden switching \([9–12]\) as well as for ac magnetic field \([13]\) was explored theoretically and experimentally for various regimes. For a constant detuning the decay of the transients is dictated by the EIT linewidth, while the frequency of the transient oscillations equals the two photon detuning \([10, 12]\). The detuning can be larger than the linewidth leading to an under-damped oscillator response. In the case of a linear sweep through the resonance the the frequency of the transient is chirped \([12]\) and behaves similarly to a Landau-Zener transition \([14, 15]\). Transients were used also as a magnetic sensing technique to measure the Earth’s magnetic field with \(1.5 \times 10^{-11}\) sensitivity \([16]\).

In this letter we take both concepts of multimode EIT and transient EIT and combine them together. Using a full mapping of the transient response as a function of the detuning we are able to show the transition from an adiabatic regime to a non-adiabatic regime. We also show the complex interference pattern that arises when a magnetic field is applied. In this case an interference between transients from three Zeeman sub-levels is visible. The problem of three level crossing in the context of Landau-Zener transitions was addressed both theoretically and experimentally \([17–22]\). Experimentally it was shown using 2 two level system (TLS) defects that have a similar energy gap inside a Josephson junction, that interference occurs at certain phases when the junction frequency is swept \([21]\). Here, on the contrary, we have the ability to change the levels energy gap experimentally implementing the three levels Landau-Zener transition for variable level detuning. Moreover, we show that this interference can be useful as a high sensitivity magnetometer combined with broadband phase modulation sweep in order to achieve a wideband high sensitivity magnetometer.

We now describe the effect a phase modulated strong coupling field has upon the temporal shape of a probe pulse going through an EIT media. The coupling field can be written as follows:

\[
E_c(t) = E_{c0}e^{i(\omega_c t - i\phi(t))},
\]

where \(E_{c0}\) is the amplitude of the coupling field, \(\omega_c\) is the optical frequency and \(\phi(t)\) is the time dependent phase modulation. In order to describe the change in the probe field due to this modulated coupling field the full spatiotemporal Maxwell-Bloch equations for an EIT system needs to be solved \([23]\). In the case of long enough interaction media and perturbative probe intensity the system can reach a steady state solution. The EIT susceptibility in this case has only temporal and no spatial dependence. For a susceptibility with no temporal dependence the probe transmission amplitude, \(p(t)\) is given by the convolution of the entering signal, \(E_p(\omega)\), and the susceptibility, \(\chi(\omega)\), hence \(p(t) = \mathcal{F}[E_p(\omega)\chi(\omega)]\). In the case of a modulated field the susceptibility is time dependent, and this convolution is not valid. A possible way of solving this problem is by taking the spectral decomposition of the susceptibility \(\chi(\omega, t) = \sum_{n=-\infty}^{\infty} e^{i\Omega n t} \chi_n(\omega)\). Now the transmission
is just \( p(t) = \sum_{n=-\infty}^{\infty} e^{i\Omega t} F'[E_p(\omega)\chi_n(\omega)] \) [23]. We use a sinusoidal modulation, hence \( \phi(t) = M \sin(\Omega t) \) where \( M \) is the modulation depth and \( \Omega = 2\pi f_c \) is the modulation frequency. The spectrum of such a modulated field can be described as a sum of Bessel functions

\[
e^{-i\phi} = \sum_{n=-\infty}^{\infty} e^{i\Omega t} J_n(-M).
\]

The spectrum of this modulation has narrow peaks separated by a frequency \( \Omega \), with a full modulation span of \( 2\Omega \). The transfer function in this case is just an infinite comb of single EIT lines [1] weighted by Bessel functions:

\[
\chi(t) = \sum_{n=-\infty}^{\infty} i\alpha \frac{J_n(-M)e^{i\Omega t}}{\Gamma - i(\Delta - n\Omega_c) + \frac{R}{\gamma_{12} - i(\delta - n\Delta)}}.
\]

Here \( \alpha = \frac{N\mu_i^2}{\epsilon_0} \) is the two-level absorption coefficient, with \( N \) the density of the atoms and \( \mu \) is the transition dipole moment, \( \Gamma \) is the homogeneous decay rate, \( \gamma_{12} \) is the decoherence rate of the two ground states, \( R = \frac{\mu_B^2}{\epsilon_0} \) is the Rabi frequency of the coupling field which is phase modulated, \( \Delta \) is the one photon detuning of the probe field and \( \delta \) is the two photon detuning.

In the case of the D1 line of warm \(^{87}\text{Rb}\) vapor with buffer gas the FWHM EIT linewidth is [24] \( \gamma_{\text{EIT}} = 2(\gamma_{12} + \frac{R}{\Gamma + |\delta|}) \) where \( \Gamma_D \) is the Doppler broadening. This linewidth is usually a few kHz which is much narrower than the pressure broadened homogeneous linewidth \( (\Gamma \sim 100 \text{ MHz}) \) and the Doppler broadening \( (\Gamma_D \sim 500 \text{ MHz}) \), thus the probe two level susceptibility is effectively constant for the full modulation bandwidth as long as \( M\Omega_c \ll \Gamma \).

Applying a magnetic field removes the Zeeman degeneracy and the energy levels of the hyperfine levels will create a ladder according to the Zeeman splitting of the two lower levels with a Larmor frequency \( \mu_B g_F m_F \) and \( g_F \). Here \( g_B \) is the magnetic field, \( \mu_B \) is the Bohr magneton and \( g_F \) is the Lande coefficient of the hyperfine level. The EIT susceptibility will be determined by the Zeeman splitting with several EIT peaks having a certain phase between them. We can write the transfer function as follows:

\[
\chi(t) = \sum_{n=-\infty}^{\infty} \sum_{m_F=-F}^{F} \sum_{m_{FF}=-F}^{F} i\alpha \frac{J_n(-M)e^{i\Omega t}}{\Gamma - i(\Delta - n\Omega_c) + \frac{R}{\gamma_{12} - i(\delta - n\Delta)}}.
\]

The experimental setup is shown in Fig. 1. For an EIT \( \Lambda \) scheme we use the hyperfine transitions of The D1 line of \(^{87}\text{Rb}\). A DFB laser locked to the \( F = 2 \rightarrow F' = 2 \) transition is split into probe and coupling beams using a polarizing beam splitter. The phase modulation over the coupling field as well as the pulse creation of the probe is done using acousto-optic modulators. In order to bring the probe to resonance with the transition an Electro-optic modulator is used. The beams (orthogonal polarization) are combined together using a Glan-Taylor polarizer and pass through a 7.5 cm cell containing an isotopically pure \(^{87}\text{Rb}\) with 10 Torr Ne as buffer gas heated to ~ 40\(^\circ\)C. The cell is shielded from an outside magnetic field using a 3-layers of \( \mu \)-metal. An axial magnetic field is created using a uniform solenoid. After the cell another polarizer is used in order to filter the coupling field while the probe is detected using an amplified photodiode.

**Figure 1.** The experimental setup. DFB - distributed feedback laser, PBS - polarizing beam splitter, AP - aperture, AOM - acousto-optic modulator, EOM - Electro-optic modulator, GT - Glan-Taylor polarizer, \( \mu \)MS - \( \mu \)-metal shield, BD - beam dump, PD - photodiode, SL - solenoid.

**Figure 2.** demonstrates the transmission temporal response of a square probe pulse with intensity of 0.05 \( \text{mW/cm}^2 \) due to a phase modulated coupling field with intensity of 1 \( \text{mW/cm}^2 \) in an EIT media. Two major features are observed, one is a train of pulses that is created with a period and phase that is dependent upon the coupling field modulation frequency and the detuning [23]. The second feature is a transient ringing that is associated with the response of the media to a sudden change in the susceptibility. This ringing has a chipped frequency as expected [12]. It decays with a characteristic time that depends upon \( 1/\gamma_{\text{EIT}} \) and the chirp rate through the transition [25]. The instantaneous frequency of the coupling field due to the modulation is \( \omega(t) = \omega_0 + \frac{\partial\alpha}{\partial t} = \omega_0 + M\Omega_c \cos(\Omega t) \) while the response has the spectral width of the EIT linewidth, thus the relation between the modulation frequency and the EIT width, sets the adiabaticity of the response.

Figure 3 shows experimentally and theoretically the transition between the adiabatic regime where the modulation frequency is lower than the EIT linewidth \( (\Omega_c \ll \gamma_{\text{EIT}}) \) and the non-adiabatic regime where \( \Omega_c \gg \gamma_{\text{EIT}} \). For both regimes the phase of the pulses is determined by the instantaneous frequency hence we see a sine like plot as a function of the detuning with a period \( 1/\Omega_c \) and an amplitude \( M\Omega_c \). In the adiabatic regime the transients decay fast enough so they are hardly noticeable, but as the modulation frequency become comparable
to the EIT linewidth [Fig. 3(b)] the transient ringing is clearly observed. In the non-adiabatic regime the modulation frequency is faster than the decay of the transient ringing creating an interference between consecutive pulses as can be observed in Fig. 3(c). Simulation of these 2D patterns using Eq. 3 are depicted in Fig. 3(c-e) showing a striking similarity to the results. One aspect this linear response theoretical simulation fails to take into account is the smearing of the interference pattern when the probe pulse is turned on as can be visualized particularly in Fig. 3(c). The cause of this effect is the gradual build up of the dark state polariton and consequently the creation of the EIT line that has a characteristic time of $\frac{1}{2\gamma_{EIT}}$ [10]. Integrating the time domain reveals the steady state spectrum of the probe light. Figure 3(f-h) shows the integrated spectra of the experimental data (green line) as well as the simulation (red line) in the adiabatic and non-adiabatic regimes. These spectra fit to the phase modulation spectrum according to the Fourier of Eq. 2, meaning a delta functions separated by the modulation frequency, broadened due to the finite EIT linewidth. Another way of thinking about the interference shown in Fig. 3(c) is as a Landau-Zener-Stuckelberg interference pattern where a transition is crossed repetitively faster than the transient decay time [26]. Figure 4 shows a 2D mapping of the temporal response of the probe for different magnetic fields (The two photon detuning is on resonance with the magnetic insensitive transition). In the adiabatic regime (Fig. 4(a)) it is possible to see a splitting of the sole pulse in B=0 into three pulses. These pulses correspond to three EIT lines that are present in the spectrum. For the D1 line of rubidium, using an arbitrary magnetic field, up to 7 EIT lines may appear [27]. Due to the vectorial nature of the magnetic interaction, the relative strength of these lines depends on the angle between the beam direction and the magnetic field as well as the polarization of the pump and probe beams [28]. The specific configuration we use in our setup is $\mathbf{B} \parallel \mathbf{k}$ with linear polarization. In this case only three lines appear in the spectrum [28, 29] as can be seen experimentally [30]. The central pulse matches the $\Delta m = 0$ and thus its phase is constant, while the other two pulses correlate with the $\Delta m = \pm 2$, hence having a sinusoidal phase shift. Figures 4(b-c) show the non-adiabatic regime where every pulse has an oscillating tail with a certain phase causing an interference pattern. A simulation based on Eq. 4 is shown in Fig. 4(d-f) having the same basic features of the experimental results.

The observed interference may be understood in the following way. In the case of hyperfine EIT in a buffer gas and under the condition of $\gamma_{EIT} \ll \Gamma$ it is possible to treat the two ground states as a degenerate set of effective two level systems [31, 32]. Adding a magnetic field removes the degeneracy and splits each two level system according to the Zeeman frequency [2]. In our case due to the selection rules stated above the splitting is to three groups with $\Delta m = 0, \pm 2$. As a consequence of this picture it is possible to think of the magnetic sweep in time as a chirp of the three TLS’s (as depicted by the black dashed lines in Fig. 4(a)). Thus, the interference we measure is a direct consequence of a three Landau-Zener transitions degeneracy [17].

The dynamic pattern created by the phase modulation can be used for broadband magnetic sensing. Each magnetic field has a certain characteristic pulse timing associated with it. The phase of the first pulse is a prominent feature for broad magnetic sensing as the total amplitude of the modulation is $B_{\text{max}} = M \Omega \gamma_{EIT}$. This corresponds to $142 mG$ and $107 mG$ for 4(b) and 4(c) respectively. Moreover, the interference pat-
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Figure 4. Temporal response of a probe pulse due to magnetic field for (a) adiabatic regime, (b) intermediate regime and (c) non-adiabatic regime. The parameters of the modulation are (a) \( f_c = 500 \text{ Hz} \) and \( M = 200 \), (b) \( f_c = 5 \text{ kHz} \) and \( M = 40 \), (c) \( f_c = 15 \text{ kHz} \) and \( M = 10 \). Plots (d-f) show a simulation of the three regimes that takes into account Eq. 4 with the parameters written above.

\[
\frac{\Delta B}{\sqrt{\Delta \nu}} = \sqrt{\frac{t}{S/N}} = \sqrt{\frac{\nu}{\partial A/\partial B}}
\]

(5)

Where \( t \) is the measurement time and \( \partial A/\partial B \) is the gradient of the integrated measured amplitude and the magnetic field. In our case, since the transient ringing is a complex multi frequency feature the best way to characterize the transients for different magnetic field is by using the correlation between them. Using this method we measured the noise and the gradient of the correlation function and estimate our sensitivity to be \( 1.2 \frac{M}{\text{VHz}} \) for the 5 kHz modulation and \( 0.2 \frac{M}{\text{VHz}} \) in the case of 15 kHz. The ultimate sensitivity for a given system is [34]

\[
\frac{\Delta B}{\sqrt{\Delta \nu}} = \frac{2}{V} \frac{\nu}{\mu_0 N V}
\]

where \( V \) is the volume of the magnetometer. In our case this sensitivity is \( \sim 400 \frac{f_c}{\text{VHz}} \) well below the measured sensitivity. The main reason for that is the electrical noise in the detector and amplifier which is used in order to observe the data in the oscilloscope. Better electronics may allow at least one order of magnitude improvement.

In conclusion, we show experimentally the transient response of an EIT media to a phase modulated pump. This response reveals explicitly the coherent nature of EIT. In the non-adiabatic regime where EIT peaks are spectrally resolved it is possible to see interference between the different modes. Albeit the interference between these modes does not contribute to EIT spectral narrowing and as a consequence to a slower light propagation [6, 35], it does create a transient behavior useful for broadband data transfer. Applying a magnetic field splits the EIT line into three, allowing us to see an interference pattern caused by the Landau-Zener crossing of the three EIT lines. Along with the wideband sweep due to the high modulation index it is shown to be a useful tool for sensitive wideband magnetometry.

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SUPPLEMENTARY MATERIAL

Landau-Zener representation

The ringing observed in Fig. 2 is a manifestation of a non-adiabatic transition through the EIT resonance. Landau-Zener theory deals with this kind of transitions and gives analytic prediction to the population transfer between the levels. In the case of an EIT in buffer gas the best way to describe the system is using the dressed state picture. Taking the Hamiltonian of the bare three levels under the rotating wave approximation

\[
\begin{pmatrix}
0 & 0 & \frac{1}{2}R_p^* \\
0 & \delta - \frac{1}{2}g_0^2 & \frac{1}{2}R_0^* \\
\frac{1}{2}R_p & \frac{1}{2}R_0 & \Delta
\end{pmatrix}
\tag{6}
\]

where \(R_0(t) = R_0 e^{-i\delta(t)}\), \(R_p\) are the Rabi frequencies of the coupling and the probe fields respectively, \(\delta\) is a constant two photon detuning and \(\Delta\) is one photon detuning where in the case relevant to us is 0. In order to see the resemblance to the Landau-Zener case it is instructive to change to a new basis where

\[
\begin{align*}
|1\rangle' &= |1\rangle \\
|2\rangle' &= |2\rangle e^{i\delta(t)} \\
|3\rangle' &= \frac{|3\rangle + e^{i\delta(t)}|2\rangle}{\sqrt{2}}
\end{align*}
\tag{7}
\]

The new Hamiltonian will become

\[
\begin{pmatrix}
0 & 0 & \frac{1}{2}R_p^* \\
0 & \delta - \frac{1}{2}g_0^2 & \frac{1}{2}R_0^* \\
\frac{1}{2}R_p & \frac{1}{2}R_0 & 0
\end{pmatrix}
\tag{8}
\]

The 2 \times 2 matrix of levels \(|2\rangle'\) and \(|3\rangle'\) is a Landau-Zener Hamiltonian. Under EIT conditions \(R_0 \gg R_p\) hence it is possible to diagonalize this 2 \times 2 matrix with two new dressed levels with eigenvalues \(\epsilon_\pm = -\frac{1}{2}\delta \pm \sqrt{\delta^2 + \frac{1}{4}g_0^2 R_0^2} + \frac{1}{2}M\Omega^2 - \frac{1}{2}R_0^2\). In the simple case where \(\delta = 0\) these states are just \(|\pm\rangle = \sin \theta |2\rangle' \pm \cos \theta |3\rangle'\) and \(|\mp\rangle = \cos \theta |2\rangle' - \sin \theta |3\rangle'\) with tan \(2\theta = \frac{2\delta R_0}{g_0^2 M\Omega^2}\). This Landau-Zener dynamics is interrogated by the probe field, meaning that the transition element \(|1\rangle \rightarrow |3\rangle\) we are measuring in the experiment, carries the dynamics described above as depicted in Fig. S5(a). In our experiment a phase modulation sweep in time causes a periodic crossing between the two dressed levels.

When a magnetic field is applied the system is split into three sub-systems with three levels in each on of them as discussed in the main page and in [2]. Each one of these sub-systems behaves exactly as a single EIT system with the exception of a magnetic Zeeman shift \(\mu_B B(g_1 m_f - g_r m_r)\). As a consequence the energy levels of the sub-systems \(\Delta m = \pm 2\) and \(\Delta m = -2\) are reversed with respect to the magnetic field (with \(\delta = 0\)) while the energy levels of the sub-systems \(\Delta m = 0\) is degenerate up to the interaction avoided level crossing as depicted in Fig. S5(b).

Figure S5. Dressed states for (a) no magnetic field (b) with magnetic field. Red arrows represent probe transition on two photon resonance (\(\delta = 0\)).

One interesting characterization of the Landau-Zener transition is the transition time. This time can be measured by the decay time of the oscillations after the transition [3]. The two parameters that determines the transition properties is the coupling Rabi frequency and the chirp rate defined as \(\frac{\partial \phi}{\partial \Omega}\). In the case of sinusoidal phase modulation, where \(\phi(t) = M\cos(\Omega t)\), the chirp rate at \(\delta = 0\) is \(\frac{\partial \phi}{\partial \Omega} = M\Omega^2\). It is useful to define the transition using a dimensionless parameter \(\beta = \frac{R_0}{\sqrt{g_0^2 M\Omega^2}}\). Figure S6 shows the decay time, \(\tau\), as a function of \(\beta\) for our experimental results (red squares) as well as for our simulation results (black circles). The decay time is found from an exponential fit to the ringing peaks as depicted in the inset in Fig. S6. It is possible to see that in the diabatic limit (low \(\beta\)) the decay time is nearly constant and converging towards \(\frac{1}{2\tau_{EIT}}\), while at the adiabatic limit (high \(\beta\)) the decay is linear with \(\beta\). Similar theoretical results for the Landau-Zener theory have been reported before [3, 4].

Temporal dynamics with magnetic field

Figure S7 shows a broad scan of magnetic field vs. time. This scan is done in the case of two photon resonance in the absence of magnetic field. We can distinct clearly the functional behavior of the three EIT lines for \(\Delta m = 0, \pm 2\). The \(\Delta m = 0\) line is not dependent upon magnetic field thus its phase is constant with a pulse every half a cycle. Both \(\Delta m = \pm 2\) lines are sinusoidally modulated with a cycle equal to \(f_1\) and a phase of \(\pi\) between them. Each of these two lines behave exactly like the detuning sweep of one EIT line (with no magnetic field) under phase modulation (see for example Fig. 3). This feature is understandable, as applying magnetic field can be translated to detuning via the Larmor frequency Zeeman shift. Adding a constant detuning or constant mag-
Magnetic field creates a symmetric shift of the two sinusoids until reaching a field larger than $2M\Omega$. In this case the two sinusoids get separated and the constant pulse of $\Delta m = 0$ disappears. Since the two sinusoids do not intersect the interference pattern disappears. The major consequence is that measuring a constant magnetic field accurately using this method is possible only for magnetic fields with Larmor frequency smaller than $2M\Omega$.

**Figure S7.** A broad scan both in time and in magnetic field sweep. Here $f_c = 2\ kHz$ and $M = 10$. (a) Experiment. (b) Simulation.

**Magnetic field splitting**

As mentioned in the main text, the spectrum of EIT under axial magnetic field creates a splitting to three sub-levels. This is certainly verified by the three pulses seen in Fig. 4(a). As a complementary measurement we also measure the steady state spectrum of the EIT under variable magnetic field as can be seen in Fig. S8.

**Figure S8.** Zeeman splitting of EIT resonance due to axial magnetic field. The observed spectrum is indeed split to three levels.

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