Electromagnetically induced transparency (EIT)
Imagine you can make a black wall transparent just by shining a certain light onto that wall. In principle, this is what happens to an atomic medium when it experiences electromagnetically induced transparency (EIT). A light beam which is completely absorbed by an atomic medium can suddenly go through just by shining in another certain light beam. A feature that raises the idea of an optical transistor, which has recently been implemented with EIT for just one atom [1].

Furthermore, the phenomenon of EIT is closely related to slow light, i.e. the reduction of the speed of light in an atomic system. A light speed reduction down to 17 m/s was achieved this way in an ultracold atomic gas [2]. Pursuing this method further, light can even be stopped, opening up the possibility of storing light in an atomic medium. If a quantum bit is encoded in that light, the atomic system can this way be made a quantum memory, as has been realized with single atoms not long ago [3]. EIT is the starting point for these developments that pave the way towards quantum communication as well as quantum computing.

The subject of this lab course is the observation and the quantitative analysis of EIT in a sample of Rubidium atoms. The preparation of EIT requires a Doppler-free saturation spectroscopy on Rubidium as well as the frequency modulation of laser beams, which is realized with acousto-optical modulators. As an application, the EIT-signal will be employed as a magnetometer to measure Bohr’s magneton.

On the practical side, the lab course includes the alignment and preparation of laser beams with opto-mechanical components as well as polarization optics, the handling of electronic devices and computer-aided analysis. On the theoretical side, the course aims at deepening the understanding of atom-light interaction, reviewing concepts like the atomic two-level system, Rabi oscillations or Zeeman splitting.

Errors, suggestions, criticism or whatever kind of feedback on this manuscript or the lab course is very welcome at EIT@matterwave.de.

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1. Introduction

Consider a three level system as indicated in figure 1.1 with three eigenstates $|1\rangle$, $|2\rangle$ and $|3\rangle$. Transitions $|1\rangle \leftrightarrow |3\rangle$ and $|2\rangle \leftrightarrow |3\rangle$ shall be allowed while the transition $|1\rangle \leftrightarrow |2\rangle$ shall be dipole-forbidden. Due to its shape, this system is called \textit{Λ-System}. It is irradiated by two laser beams named \textit{probe beam} and \textit{couple beam} with angular frequencies $\omega_p$ and $\omega_c$, respectively.

![Figure 1.1: Schematic of a general Λ-System. The couple beam with angular frequency $\omega_c$ couples eigenstates $|2\rangle$ and $|3\rangle$ while the probe beam with angular frequency $\omega_p$ probes the transition $|1\rangle \leftrightarrow |3\rangle$.](image)

If the couple beam is switched off while the frequency of the probe beam is varied one would expect an absorption profile around the resonance frequency, as illustrated on the top of figure 1.2.

![Figure 1.2: First observation of EIT in Strontium vapor](image)

When the couple beam is switched on, a tiny transparency window occurs in the absorption valley of the otherwise opaque atomic medium.
If the couple beam is turned on a peculiar phenomenon occurs: A tiny transparency window opens around the resonance frequency, reaching a transmission of nearly 50% in an otherwise opaque surrounding (bottom of figure 1.2). This effect is called electromagnetically induced transparency (EIT) and was first observed in Strontium vapor by [4], whose results are shown in figure 1.2.

Simply speaking, this means that a medium can be made transparent for a certain frequency by switching on a couple beam.

In the next chapter an attempt will be made to provide a theoretical understanding of this not very intuitive EIT-phenomenon. On top of that, the equations for calculating the EIT-features are deduced.
2. Theory of EIT

To obtain a theoretical understanding of EIT, a quantum mechanical treatment of the three-level system is necessary. Such a treatment builds on the theory of the two-level system, which should have been subject of the quantum mechanics lecture. For memorization a brief recalling of the two-level system will be given in the following section, where the underlying concepts are introduced. Based on that, the theory of EIT will first be described in a static way only looking at the eigenstates of the three-level system. For obtaining the dynamic features of EIT as observed in the experimental course, the three-level system will then be treated dynamically.

2.1. The two-level atom

A detailed derivation of the two-level system can be found in for example [5, 6, 7, 8, 9]. The approach presented here mainly follows [5] and [9]. Consider an atom that has only a ground state $|g\rangle$ and an excited state $|e\rangle$. These are eigenstates of the atomic Hamiltonian $\mathcal{H}_0$, i.e.

$$\mathcal{H}_0 |e\rangle = \hbar \omega_e |e\rangle$$

$$\mathcal{H}_0 |g\rangle = \hbar \omega_g |g\rangle$$

The atom shall now be illuminated with laser light of the angular frequency $\omega_l$ corresponding to a detuning $\delta = \omega_l - \omega$ with $\omega = \omega_e - \omega_g$.

![Figure 2.1.: Scheme of the idealized two-level atom with ground state $|g\rangle$ and excited state $|e\rangle$ that have energies $\hbar \omega_g$ and $\hbar \omega_e$ respectively. The two-level system is irradiated by light with angular frequency $\omega_l$.](image)

Assuming a classical light field, the time-evolution of the system can be described with the Schrödinger equation. However, including spontaneous emission requires to describe the system by the density matrix

$$\rho = \begin{pmatrix} \rho_{ee} & \rho_{eg} \\ \rho_{ge} & \rho_{gg} \end{pmatrix}$$

(2.1)
where $\rho_{ee}$ and $\rho_{gg}$ are the probabilities to find the system in state $|e\rangle$ and $|g\rangle$, respectively. The off-diagonals $\rho_{eg}$ and $\rho_{ge}$ determine the coherence between the two states. The time evolution of the density and thereby the time evolution of the state populations is obtained by the von Neuman equation

$$i\hbar \dot{\rho} = [\mathcal{H}, \rho]$$

(2.2)

where $\mathcal{H} = \mathcal{H}_0 + \mathcal{H}'$ denotes the Hamilton operator which is composed of the atomic Hamiltonian $\mathcal{H}_0$ and the Hamiltonian $\mathcal{H}'$ that describes the atom–light interaction. In the dipole approximation this Hamiltonian can be expressed as

$$\mathcal{H}' = -d \cdot E$$

(2.3)

where $d$ is the electric dipole operator and $E$ the electric field produced by the laser at the location of the atom. Evaluating the differential equation (2.2) and introducing the Rabi-frequency

$$\Omega := \frac{d \cdot E}{\hbar}$$

(2.4)

yields the so-called Optical Bloch Equations (OBE)

$$\begin{align*}
\dot{\rho}_{gg} &= \gamma \rho_{ee} + \frac{i}{2} (\Omega^* \rho_{eg} - \Omega \rho_{ge}) \\
\dot{\rho}_{ee} &= -\gamma \rho_{ee} + \frac{i}{2} (\Omega \rho_{ge} - \Omega^* \rho_{eg}) \\
\dot{\rho}_{ge} &= -\frac{\gamma}{2} \rho_{ge} + \frac{i}{2} [\Omega (\rho_{ee} - \rho_{gg}) - 2\delta \rho_{ge}] \\
\dot{\rho}_{eg} &= -\frac{\gamma}{2} \rho_{eg} + \frac{i}{2} [\Omega^* (\rho_{gg} - \rho_{ee}) + 2\delta \rho_{eg}]
\end{align*}$$

(2.5)

where the first summand has been added to account for spontaneous emission\footnote{this is a somewhat sloppy way of introducing the spontaneous emission. However, a more rigorous derivation\cite{10} leads to the same result, but would exceed the scope of this manual.} that occurs at the spontaneous decay rate $\gamma$. Furthermore, $\tilde{\rho}_{ge} = \rho_{ge} e^{i\delta t}$ and $\tilde{\rho}_{eg} = \rho_{eg} e^{i\delta t}$ have been introduced and $^*$ denotes the complex conjugate.

The spontaneous decay rate $\gamma$ is obtained from a fully quantized theory\footnote{The correct expression for the spontaneous decay rate was already found by a phenomenological theory of Einstein\cite{5}.}, where also the electromagnetic field is treated quantum mechanically\footnote{This is a somewhat sloppy way of introducing the spontaneous emission. However, a more rigorous derivation\cite{10} leads to the same result, but would exceed the scope of this manual.}. It is given by

$$\gamma = \frac{\omega^3 e^2 d^2}{3\pi \epsilon_0 \hbar c^3}$$

(2.6)

Using conservation of population

$$\rho_{gg} + \rho_{ee} = 1$$

(2.7)

and optical coherence

$$\rho_{eg} = \rho_{ge}^*$$

(2.8)

these equations can further be reduced to

$$\begin{align*}
\dot{\rho}_{eg} &= -\left(\frac{\gamma}{2} - i\delta\right) \rho_{eg} + \frac{i\omega \Omega}{2} \\
\dot{w} &= -\gamma w - i \left(\Omega^* \rho_{eg} - \Omega \rho_{eg}^*\right) + \gamma
\end{align*}$$

where $w := \rho_{gg} - \rho_{ee}$ denotes the difference in population.
The temporal behaviour of the system can be found by direct numerical integration of the optical Bloch equations. The result for different detunings and Rabi-frequencies is shown in figure 2.2.

Figure 2.2.: Rabi oscillations for different detunings without spontaneous emission (left) and including spontaneous emission for different Rabi-frequencies (right).

Without spontaneous emission (i.e. $\gamma = 0$) the two-level system performs oscillations between the ground and the excited state. In case of resonance ($\delta = 0$) these oscillations are fully modulated and occur at the Rabi-frequency $\Omega$, whereas for $\delta \neq 0$ the modulation of the oscillations decreases and the frequency increases to

$$\Omega' = \sqrt{\delta^2 + \Omega^2}.$$  \hfill (2.9)

If spontaneous emission is included (i.e. $\gamma \neq 0$) one can distinguish between two different time domains:

- For $t \ll \frac{1}{\gamma}$ the damping caused by spontaneous emission is negligible.
- For $t \gg \frac{1}{\gamma}$ a steady-state will be reached due to damping caused by spontaneous emission.

The steady-state population of the excited state $\rho_{ee}$ in the second case can be obtained by requiring

$$\dot{\rho}_{eg} = \dot{w} = 0$$  \hfill (2.10)

which yields

$$\rho_{ee} = \frac{s/2}{1 + s + \frac{4\delta^2}{\gamma^2}}$$  \hfill (2.11)

where the saturation parameter

$$s := \frac{2\Omega^2}{\gamma^2} = \frac{2|d|^2|E|^2}{\hbar\gamma^2} = \frac{I}{I_s}$$  \hfill (2.12)

has been defined. Furthermore, the saturation intensity

$$I_s := \frac{\hbar c \epsilon_0 \gamma^2}{4d^2} \frac{2\pi h c \gamma}{3\lambda^3}$$  \hfill (2.13)

is introduced, while using that the intensity of the light field is related to the electric field via

$$I = \frac{1}{2} c \epsilon_0 |E|^2.$$  \hfill (2.14)

For the case of low saturation with $s \ll 1$, the population is predominantly in the ground state, whereas in the case of high saturation, the population is equally distributed between ground and
excited state.

From (2.11) the total scattering rate $\gamma_s$ of light from the laser field is given by

$$\gamma_s = \gamma \rho_{ee} = \frac{\gamma s}{2(1 + s + \frac{4s^2}{\gamma^2})}$$  \hspace{1cm} (2.15)

which saturates for large $s$ to the maximum scattering rate $\gamma/2$.

Plotting the scattering rate as a function of detuning $\delta$ (figure 2.3) yields a Lorentzian-profile with a FWHM of

$$\gamma' = \gamma \sqrt{1 + s}$$  \hspace{1cm} (2.16)

which is referred to as the power-broadened linewidth.

![Figure 2.3: The scattering rate $\gamma_s$ as a function of detuning $\delta$ for different saturation parameter $s$ is shown. For $s > 1$ the line profiles are substantially broadened.](image)

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2.2. Static description of EIT

Having worked through the two-level system, the three-level system can be approached. For a theoretical description of EIT, the simple lambda system in figure 1.1 shall be elaborated a bit further. To that objective the relevant parameters are added in figure 2.4.

![Diagram of a Lambda System with Rabi frequencies and detunings](image)

Figure 2.4.: Schematic of a Λ-System with Rabi frequencies $\Omega_p$ and $\Omega_c$ driven by probe and couple beam respectively, spontaneous decay rates $\gamma_{ij}$ for the decay from level $|i\rangle$ to $|j\rangle$ and the detunings $\Delta_p$ and $\Delta_c$ of the beams relative to the corresponding transition.

$\Omega_p$ and $\Omega_c$ denote the Rabi frequencies of the transitions driven by the probe-beam and the couple-beam respectively, which are related to the corresponding beam intensities $I_p$ and $I_c$ by

$$\Omega_p^2 = \frac{3\gamma_{31}\lambda^3}{4\pi^2\hbar c} \cdot I_p \quad \Omega_c^2 = \frac{3\gamma_{32}\lambda^3}{4\pi^2\hbar c} \cdot I_c \quad (2.17)$$

The detunings of the probe and the couple frequency relative to the corresponding transition frequencies are

$$\Delta_p = \omega_{31} - \omega_p \quad \Delta_c = \omega_{32} - \omega_c \quad (2.18)$$

where $\omega_{31} = \omega_3 - \omega_1$ and $\omega_{32} = \omega_3 - \omega_2$ denote the angular frequency differences between the corresponding eigenstates. The spontaneous decay rates $\gamma_{31}$ and $\gamma_{32}$ account for the spontaneous decay from state $|3\rangle$ to $|1\rangle$ and $|2\rangle$ respectively. Although the transition from $|2\rangle$ to $|1\rangle$ is dipole forbidden there still might be a tiny probability for a spontaneous decay accounted for by $\gamma_{21}$.

The crucial point for EIT is that the presence of the probe and couple beam changes the Hamiltonian, such that an interaction Hamiltonian $H_{\text{int}}$ describing the interaction between the atom and the incoming light has to be added to the mere atomic Hamiltonian $H_0$:

$$H = H_0 + H_{\text{int}} \quad (2.19)$$

The eigenstates of $H_0$, $|1\rangle$, $|2\rangle$ and $|3\rangle$, are in general no longer eigenstates of the total Hamiltonian $H$. After having introduced the dipole approximation as well as the rotating wave approximation, the interaction Hamiltonian can be represented in a rotating frame by

$$H_{\text{int}} = -\frac{\hbar}{2} \begin{pmatrix}
0 & 0 & \Omega_p \\
0 & -2(\Delta_p - \Delta_c) & \Omega_c \\
\Omega_p & \Omega_c & -2\Delta_p
\end{pmatrix} \quad (2.20)$$
The eigenvalues of the interaction Hamiltonian $\mathcal{H}_{\text{int}}$ can be calculated to

$$\hbar \omega^+ = \frac{\hbar}{2} \left( \Delta_p + \sqrt{\Delta_p^2 + \Omega_p^2 + \Omega_c^2} \right)$$  \hfill (2.21)

$$\hbar \omega^- = \frac{\hbar}{2} \left( \Delta_p - \sqrt{\Delta_p^2 + \Omega_p^2 + \Omega_c^2} \right)$$  \hfill (2.22)

$$\hbar \omega^0 = 0 .$$  \hfill (2.23)

It is straightforward to verify that the corresponding eigenstates of the interaction Hamiltonian $\mathcal{H}_{\text{int}}$ are

$$|a^+\rangle = \sin \theta \sin \phi |1\rangle + \cos \theta \sin \phi |2\rangle + \cos \phi |3\rangle$$  \hfill (2.24)

$$|a^-\rangle = \sin \theta \cos \phi |1\rangle + \cos \theta \cos \phi |2\rangle - \sin \phi |3\rangle$$  \hfill (2.25)

$$|a^0\rangle = \cos \theta |1\rangle - \sin \theta |2\rangle$$  \hfill (2.26)

where the mixing angles $\theta$ and $\phi$ with

$$\tan \theta = \frac{\Omega_p}{\Omega_c} \quad \text{and} \quad \tan \phi = \frac{\sqrt{\Omega_p^2 + \Omega_c^2}}{\sqrt{\Omega_p^2 + \Omega_c^2 + \Delta_p^2}}$$  \hfill (2.27)

have been introduced. Now the state $|a^0\rangle$ shall be considered in detail, especially the question is asked: how does $|a^0\rangle$ couple to state $|3\rangle$ or in other words, what is the transition probability to go from state $|a^0\rangle$ to state $|3\rangle$ when $\mathcal{H}_{\text{int}}$ is applied?

$$\langle 3 | \mathcal{H}_{\text{int}} | a^0 \rangle = 0$$  \hfill (2.28)

This means, that $|a^0\rangle$ does not interact with the light or in other words, a transition from $|a^0\rangle$ to $|3\rangle$ is not possible. This is why $|a^0\rangle$ is called a dark state. This provides a formal understanding of the electromagnetically-induced transparency described in the previous section: the couple-beam transfers the atom into a dark state from which an interaction with the probe light is no longer possible. Therefore, it appears to be transparent.

Now consider the special case that $\Omega_p \ll \Omega_c$, i.e. the couple beam is much stronger than the probe beam. In this case

$$\sin \theta \rightarrow 0 \quad \cos \theta \rightarrow 1 \quad |a^0\rangle \rightarrow |1\rangle$$  \hfill (2.29)

So the original state $|1\rangle$ becomes the dark state.
2.3. Dynamic description of EIT

With the static description above, the basic phenomenon of EIT could be explained. In order to derive the experimentally observed absorption profile, a dynamic description is needed that also includes spontaneous emission (note that spontaneous emission has been neglected in the previous treatment). As for the two-level system, this can be done with the density matrix formalism. Again, $\rho$ denotes the density matrix, whose diagonal elements $\rho_{ii}$ represent the measurement probability of state $i$, while the off-diagonal elements $\rho_{ij}$ are referred to as *coherences*, since they provide information about the relative phase of state $i$ and $j$.

The time evolution of the density matrix is given by

$$ i\hbar \dot{\rho} = [H_{\text{int}}, \rho] . $$

(2.30)

This equation does not account for spontaneous emission yet. Therefore decay terms have to be added which occur at rates $\gamma_{31}$ and $\gamma_{32}$ for spontaneous emissions $|3\rangle \rightarrow |1\rangle$ and $|3\rangle \rightarrow |2\rangle$, respectively. This procedure yields a set of differential equations that can be solved for each element of the density matrix $\rho$.

The final objective is to derive an expression for the transmission $T$ which is given by *Lambert-Beers Law*

$$ T(\Delta_p) = e^{-a(\Delta_p)L} $$

(2.31)

where $a$ denotes the absorption coefficient and $L$ the length of the atomic sample. The absorption coefficient can directly be calculated from $\rho_{13}$ which in turn is obtained by solving the above differential equation. For a sufficiently weak probe field ($\Omega_p \ll \gamma_{31}$) and $\Delta_c = 0$, the absorption coefficient can be approximated by

$$ a(\Delta_p) = \frac{n \omega_p \mu_{13}^2}{\epsilon_0 c \hbar} \cdot \frac{2\gamma_{21}(\gamma_{21}\gamma_{31} + \Omega_p^2) + 8\Delta_p^2 \gamma_{31}}{(4\Delta_p^2 - \gamma_{21}\gamma_{31} - \Omega_p^2)^2 + 4\Delta_p^2(\gamma_{21} + \gamma_{31})^2} $$

(2.32)

absorption coefficient

where $n$ denotes the atomic number density and $\mu_{13}$ the transition matrix dipole element. Assume that a probe beam shines through an atomic sample while scanning its frequency, i.e. its detuning $\Delta_p$ relative to the resonance frequency $\omega_{31}$. The simplification shall be made, that all atoms have zero velocity, such that the doppler broadening can be neglected. In this case the resulting transmission, i.e. the signal one would expect on a photodiode behind the atomic sample can be calculated with the above formula (2.32). If the couple beam is turned off, i.e. $\Omega_c = 0$, and if $\gamma_{21} = 0$, i.e. there is no "leakage" from state $|2\rangle$ to state $|1\rangle$, then (2.32) reduces to

$$ a(\Delta_p) = \frac{n \omega_p \mu_{13}^2}{\epsilon_0 c \hbar} \cdot \frac{2\gamma_{31}}{4\Delta_p^2 + \gamma_{31}^2} $$

(2.33)

which is the well-known natural line profile with a full-width-at-half-maximum (FWHM) of $\gamma_{31}$. The resulting transmission $T(\Delta_p)$ is plotted on the left in figure 2.3. If now the couple beam is switched on, i.e. $\Omega_c \neq 0$ then a narrow transparency window emerges at $\Delta_p = 0$ (right hand side of figure 2.3), which very much resembles the experimental result shown in the introduction. In [12] it is shown, that for $\Omega_p, \Omega_c \ll \gamma_{31}, \gamma_{32}$ this profile is approximately the difference of the above broad Lorentzian and a narrow Lorentzian with a FWHM of

$$ \gamma_{\text{EIT}} = \frac{\Omega_p^2 + \Omega_c^2}{\gamma_{31} + \gamma_{32}}. $$

(2.34)

This implies the astonishing fact, that the linewidth of the EIT-signal can in principle be made arbitrarily small, far below the natural linewidth, just by decreasing the intensity of the probe and the couple beam.
Figure 2.5.: Transmission calculated with formula (2.32) with and without couple beam switched on. \( \gamma_{21} = 0.02 \cdot \gamma_{31} \) for both plots. With the applied couple beam a narrow transmission window emerges at resonance frequency (right), which is the EIT signal.

For \( \Omega_c \gg \gamma_{13} \) the transparency window becomes very large and two Lorentz-Profiles emerge that are split by the Rabi-frequency \( \Omega_c \). This phenomenon is referred to as \textit{Autler-Townes splitting} and is conceptually different from EIT [13].

As mentioned, the theoretical description was so far only valid for atoms with zero velocity, which is nearly the case in an ultracold atomic gas. In a hot gas, as used in this experiment, Doppler-broadening has to be accounted for. However, it turns out that for co-propagating couple and probe beam, narrow EIT-resonances can also be obtained in a hot gas [14].

### 2.4. Experimental realization of an EIT-system

So far the description of EIT was made for a general \( \Lambda \)-system, which for an experimental realization needs to be prepared in an atomic system. In this experiment the D1-line of \(^{87}\text{Rb} \) is employed as indicated in figure 2.4.

\(^{87}\text{Rb} \) has a nuclear spin of \( I = 3/2 \), therefore the lower \( 5^2 S_{1/2} \) and the upper \( 5^2 P_{1/2} \) each split into hyperfine-levels with \( F = I \pm J = \{1, 2\} \). Each of these levels consists of \( 2F + 1 \) Zeeman-sublevels (which are degenerate in the absence of a magnetic field), as illustrated in figure 2.4 for the case of the lower \( F = 2 \) and the upper \( F = 1 \) level.

The question is, how a \( \Lambda \)-system can be realized in that particular atomic setting. One \( \Lambda \)-system is for example realized by the three states marked in red. Since selection rules require

\[ \Delta M_F = \pm 1 \text{ or } 0 \quad (2.35) \]

for the magnetic quantum number, the \( |1\rangle \leftrightarrow |2\rangle \) transition is dipole forbidden, thus fulfilling the prerequisite for a \( \Lambda \)-system. Without the presence of a magnetic field, couple and probe beam have the same frequency, but different polarizations recalling that \( \Delta M_F = +1 \) transitions are induced by \( \sigma^+ \) light while \( \Delta M_F = -1 \) transitions are induced by \( \sigma^- \) light. When a magnetic
field of strength $B$ is applied, the magnetic sublevels split, i.e. the energy of a magnetic substate with quantum number $M_F$ is shifted by the frequency
\[
\Delta \omega_Z = M_F g_F \mu_B / \hbar
\]
as indicated on the right in figure 2.4. Due to this relation, EIT can serve as a magnetic field sensor and in this experiment, Bohrs magneton will be measured this way.

In the following some optical transition properties of the $^{87}$Rb D1-line are listed:

<table>
<thead>
<tr>
<th>wavelength</th>
<th>$\lambda$</th>
<th>$\gamma = \gamma_{31} = \gamma_{32}$</th>
<th>$\gamma_{12}$</th>
</tr>
</thead>
</table>
| decay rate /Natural line width (FWHM) | $\gamma_{31} = \gamma_{32}$ | $2\pi \cdot 5.746(8)$ MHz | $2.537 \cdot 10^{-29}$ C \cdot m
| transition dipole moment | $\mu$ | $2.537 \cdot 10^{-29}$ C \cdot m |

These values can be inserted into expression 2.32 using typical densities ranging from $n = 10^{16} \ldots 10^{22}$ atoms $m^{-3}$. It is recommended to plot the resulting transmission for different values of $\Omega_c$, $\gamma_{12}$ and $n$ to get a feeling for the effect of the different parameters. Since the transition $|1\rangle \Leftrightarrow |2\rangle$ is dipole forbidden $\gamma_{12} \ll \gamma_{31}$ can be chosen or $\gamma_{12}$ may even be set to zero.
3. EIT-related phenomena

EIT is the starting point for many further physics. As will be revealed during the practical course, the narrow EIT-signal can be used as a magnetometer for measuring low magnetic fields with high precision ([15, 16, 17]). The property of having the control to let a probe beam pass or not by switching on a couple beam raises the idea of an optical transistor. This technique has been scaled down by Muecke et al. ([1]) to just one single atom in a cavity acting as a quantum-optical transistor for single photons. In this section a brief glance is shed on how the EIT-resonance can be used for realizing *slow light* and *light storage*.

3.1. Slow light

In the theory section the focus laid on the absorptive features of EIT and therefore an expression for the absorption coefficient $a$ was aimed at. It is related to the imaginary part of the refractive index $n$ by

$$a = I(n) .$$  

(3.1)

The real part of the refractive index instead describes the dispersive properties, i.e. how light propagates through the medium. The group velocity of the light traveling through the medium

![Figure 3.1.: Real and imaginary part for the refractive index as a function of angular frequency. At resonance where the absorption is strongest, the derivative of the real part of the refractive index also obtains its maximum, which leads to a small group velocity $v_g$ in the medium.](image-url)
is related to the real part of the refractive index $R(n)$ by:

$$v_g \propto \left( \frac{dR(n)}{d\omega} \right)^{-1} \quad (3.2)$$

i.e. the group velocity is inversely proportional to the derivative of $R(n)$ with respect to the angular frequency $\omega$. This relation implies that the steeper the dispersion curve in a medium, the slower the light will propagate. Figure 3.1 shows the real and the imaginary part of the refractive index at an EIT-resonance. The steepness of the dispersion curve is related to the width of the EIT signal, i.e. the smaller the EIT-width the steeper the dispersion curve and the slower the light. More general it can be shown that

$$v_g \propto \frac{\Omega_c^2}{N} \quad (3.3)$$

which implies that slower velocities can be achieved by either decreasing the intensity of the couple beam (i.e. reducing the EIT-linewidth) or by increasing the atomic density $N$. Exactly this is how Hau et al. [2] achieved to reduce the speed of light down to $17 \frac{\text{m}}{\text{s}}$ (figure 3.2), by going to high atomic densities in an ultracold atomic gas, while inducing EIT with a couple beam of low intensity. They observed a time delay in a light pulse propagating through an atomic sodium sample, which at lowest temperature corresponds to a group velocity of $v_g \sim 17 \frac{\text{m}}{\text{s}}$.

Figure 3.2.: Slow light in an ultracold gas of sodium atoms observed by Hau et al. [2]. The ingoing reference pulse (open circles) is delayed by $\sim 7 \mu\text{s}$ while traveling through the $\sim 229 \mu\text{m}$ long atomic cloud. The corresponding light speed is $32.5 \frac{\text{m}}{\text{s}}$. 
3.2. Light storage

The phenomenon of slow light raises the question, whether it is possible to make the light stop entirely. Recalling equation (3.3) for the group velocity it results, that upon adiabatically ramping down the couple beam power while the probe pulse is inside the atomic sample, the group velocity can actually be brought down to zero. The probe pulse remains spatially compressed inside the atomic cloud until the couple beam is switched on again. This effect was first observed by Liu et al. [18], whose results are shown in figure 3.3.

Figure 3.3.: Light storage observed by Liu et al. [18]. In a) the reference pulse (open circles) enters the atomic sample while the couple beam (dashed curve) is switched on. A delay of the outgoing pulse is observed as explained in the previous section. In b) the couple beam is adiabatically switched off when the incoming probe pulse is entirely enclosed in the atomic sample. Upon switching on the couple beam the light pulse comes out again. Storage times of more than a second were achieved this way.

Specht et al. even achieved to realize this method for coherently storing a single photon in a single atom [3]. This way they were able to implement a single-atom quantum memory, by performing write and read operations of arbitrary polarization states of light into and out of a single atom trapped inside an optical cavity.
4. Experimental setup

The aim of this advanced lab course is the observation and the quantitative analysis of EIT. The \( \Lambda \)-system is generated in the D1-line of \( ^{87}\text{Rb} \), as described in section 2.4. In the following the experimental setup shown in figure 4.2 will be explained. The employed devices are described in detail in the following chapter. The experimental setup can be divided into the following 4 parts:

![Figure 4.1.: Impression of the experimental setup in the laboratory](image)

![Figure 4.2.: Schematic of the Experimental setup for the observation of EIT](image)
4.1. Laser beam generation

A DFB-diode laser provides a light beam at 795 nm which passes an optical isolator, that protects the laser from back-reflection which can cause frequency instabilities and in the worst case damage. With the subsequent beam shaper the light beam, which is leaving the laser with an elliptical shape, is made circular. The beam shaper is a polarization dependent device, which is why a \( \lambda/2 \)-plate is placed in front for aligning the polarization. This part of the experiment is already aligned and must not be touched.

4.2. Spectroscopy

As described in section 2.4, the D1-line of \(^{87}\text{Rb}\) will be employed for EIT. Therefore the laser has to be adjusted to the frequency of that line. In order to find it, Doppler-free saturation spectroscopy will be performed. For that reason, a fraction of the laser beam is diverted by the \( \lambda/2 \)-plate and the PBS, together acting as a power splitter. The laser beam which has passed the Rubidium sample twice than goes undeflected through the PBS, since it went through the \( \lambda/4 \)-plate twice. Behind the PBS it is diverted onto a photodiode. By scanning the laser frequency (for example by scanning the lasercurrent) a Doppler-free spectrum of the D1-line can be observed and the desired frequency positioned.

4.3. Frequency generation

For EIT a probe and a couple beam are required. As explained in section 2.4 they need to have opposite polarizations and in the non-degenerated case (i.e. when a magnetic field is applied) also their frequencies are different. Furthermore, the frequency of the probe beam has to be scanned in order to observe the desired transparency profile.

A convenient optical device standardly used to shift the frequency of a laser beam is an acousto-optical-modulator (AOM). Here, AOM’s in double-pass configuration are used to adjust the frequencies of couple and probe beam. The frequency-tuned beams with orthogonal polarization are subsequently combined in a polarizing beam splitter and finally coupled into an optical fiber. Optimal coupling into the fiber guarantees a perfect overlap of couple and probe beam as well as a Gaussian beam leaving the other end of the fiber.

4.4. EIT-cell

The overlapping probe and couple beam, that are leaving the fiber have to be guided through the EIT-cell. The EIT-cell is similar to the spectroscopy cell, but optimized for optimal EIT performance. In order to increase the atomic density the Rubidium (which to 72% consists of \(^{85}\text{Rb}\)) is enriched with \(^{87}\text{Rb}\). Furthermore, the cell walls can be heated such that the rubidium atoms do not condense there, which would lower the atomic density. Suppression of disturbing magnetic fields is achieved by several layers of Mu-metal (a metal with high permeability) as can be seen in figure 4.3. The Mu-metal can be demagnetized by a degaussing-technique with coils that are wound around the Mu-metal layers. Controlled magnetic fields can be applied to the cell via a coil, which is placed inside of the magnetic shielding. A \( \lambda/4 \)-plate in front of the cell changes the linear polarization of the probe and the couple beam to circular polarization.
Behind the cell the polarization is again changed back to linear polarization such that the probe beam passes through the adjacent PBS while the couple beam is deflected. For a very efficient separation of probe and couple beam, a Glen-Taylor polarizer is used. The probe beam is then focused onto a photodiode, where the EIT-signal can be observed.

Figure 4.3.: Opened EIT-cell. A glass cell very similar to the spectroscopy cell is filled with enriched $^{87}$Rb. The cell can be heated with resistance wires in order to increase the atomic density. The magnetic field is generated with a coil of 98 mm length, 41 mm diameter and 167 windings. The cell is enclosed in several layers of Mu-metal for magnetic shielding. This shielding has to be demagnetized from time to time by applying deegaussing-currents to wires wound around the Mu-metal layers.
5. Devices in detail

In this section the devices employed in the experimental setup are described in more detail. However, only the properties relevant for this experiment will be explained. For a more detailed discussion for example [19] or [20] may be consulted.

5.1. $\frac{\lambda}{2}$-plate and $\frac{\lambda}{4}$-plate

A $\frac{\lambda}{2}$-plate is a birefringent crystal, which is cut such that the optical (also called extraordinary or fast) axis is parallel to the entrance surface. Light polarized along this axis travels faster than light polarized along the perpendicular axis in the plane of the surface (slow or ordinary axis). Depending on the thickness of the crystal, light with polarization components along both axes will emerge in a different polarization state. If the thickness is chosen, such that the crystal causes a phase shift between the two components of $\lambda/2$, the linear polarization is turned by $90^\circ$ (top of right hand figure).

If the thickness of the crystal is chosen, such that it causes a phase shift of $\lambda/4$, linearly polarized light is changed to elliptically polarized light and vice versa. In the special case that the angle between optical axis and incoming linear polarization is $45^\circ$, linear polarization is changed to circular polarization (bottom of right hand figure).

5.2. Polarizing Beamsplitter (PBS)

A polarizing beamsplitter (PBS) is used to separate or combine light with perpendicular polarization components. The most common type consists of two prism that are glued together with a dielectric coating in between, that has a high reflectivity for vertically (or $s$-) polarized light and a low reflectivity for horizontally (or $p$-) polarized light. Typically the transmission of the $s$-polarized light $T_s$ is suppressed by a factor of $\sim 1000$ with respect to the transmission of the $p$-polarized light $T_p$ (i.e. $T_p/T_s \sim 1000$), whereas for the reflected beam its only $T_s/T_p \sim 100$. Due to the very thin layer of dielectric coating and the cubic design, the PBS only causes a negligible shift to the path of the transmitted beam.

A PBS together with a $\lambda/2$-plate in front is standardly used as a powersplitter.
5.3. Glan-Taylor polarizer

If a purer separation of the polarization is needed a Glan-Taylor polarizer can be employed, which features an extinction ratio of \( T_p/T_s \sim 100.000 \) for the transmitted beam (the reflected beam is not that well polarized). In opposite to the PBS described above, the separation of the polarization components is the result of birefringency and total internal reflection. The Glan-Taylor polarizer consists of two birefringent prisms that are separated by a tiny air slit (\( \sim 50 \mu m \)). The optical axes of the two prisms lie horizontally in the plane of the front surface. Hence, the \( p \)-polarized (extraordinary) beam and the \( s \)-polarized (ordinary) beam are traveling along the same path, but with different velocities due the different refractive indices the two beams experience. The cutting angle of the two prisms is chosen according to the different refractive indices for \( s \)- and \( p \)-polarization, such that the \( s \)-polarized light experiences total internal reflection, while for the \( p \)-polarized component the angle is close to the Brewster angle. The ordinary ray is therefore reflected, while the extraordinary ray is transmitted. Since the air gap is very tiny, the path of the transmitted beam remains nearly unshifted. The reflected beam instead leaves the cube with an angle typically around 70°.

5.4. DFB laser diode

In this experiment a DFB (Distributed feedback) laser diode at 795nm is used. In usual Fabry-Perot diodes (top of right hand figure) the polished front surfaces of the diode build a cavity for frequency selection. Its linewidth is typically on the order of several nm. A DFB diode instead has a diffraction grating within the active region as a frequency selective element, with which a linewidth of only 1MHz to 5MHz is achieved while being able to scan over a large frequency range (several 100 GHz). The frequency of the DFB-laser can be varied by either changing the temperature or the current. In order to keep the frequency constant, the temperature of the diode is stabilized by a peltier element.

5.5. Optical isolator

The optical isolator is used to prevent backreflections into the laser, which can cause frequency instabilities and in the worst case damage. It consists of a polarizer (for example a PBS) that only lets through a certain polarization followed by a faraday rotator that changes the polarization by 45° through the faraday effect. The subsequent polarizer is also rotated by 45°, such that light traveling in this direction...
remains unaffected. However, light entering backwards will be polarized by the analyzer, turned to horizontal polarization by the Faraday rotator and then blocked by the polarizer. Hence, light can only pass the optical isolator in one direction.

5.6. Anamorphic prism pair

An elliptical beam can be shaped with an anamorphic prism pair, i.e. one axis can be expanded or compressed while the other axis stays unchanged. The prisms are usually oriented, such that they are near Brewster condition to ensure maximum transmission. Therefore a $\lambda/2$-plate for aligning the polarization is placed in front.

5.7. Doppler-free saturation spectroscopy

If a laser is sent through an atomic sample with its laser frequency scanned around the atomic resonance frequency, the transmission signal behind is a Doppler-profile, which is caused by the velocity distribution of the atoms (top of figure 5.1). The Doppler-width is typically on the order of $\sim 1\,\text{GHz}$ which is much larger than the natural linewidth (typically on the order of $\sim 10\,\text{MHz}$). In order to resolve the natural line profile, Doppler-free saturation spectroscopy can be performed.

![Figure 5.1: Doppler-free saturation spectroscopy](image)

Figure 5.1.: Doppler-free saturation spectroscopy. In the setup shown on the top only the Doppler-broadened absorption profile can be observed on a photodiode. By passing the spectroscopy twice with counterpropagating laser beams, as illustrated on the bottom, the natural line profile can be resolved.

Its idea is illustrated in figure 5.2. Without excitation, the atoms are populated in the ground state according to a Maxwell-Boltzmann velocity distribution which in one dimension is a Gaussian curve. An atom with velocity $\vec{v}$ encountering a laser beam with wavevector $\vec{k}$ experiences a
laser frequency shifted by the Doppler-shift

$$\Delta f_D = -\frac{1}{2\pi} \vec{k} \cdot \vec{v}$$  \hfill (5.1)

If the laser frequency \(f_l\) is detuned from the atomic resonance frequency \(f_0\) by \(\Delta f = f_l - f_0\), then the frequency experienced by the atom is

$$f_0 + \Delta f - \Delta f_D.$$  \hfill (5.2)

If the detuning matches the Doppler-shift, i.e. if

$$\Delta f = -\frac{1}{2\pi} \vec{k} \cdot \vec{v}$$  \hfill (5.3)

then the atom is in resonance. If the detuning is negative, an atom traveling antiparallel to the laser beam (i.e. \(\vec{k} \cdot \vec{v} < 0\)) with a certain velocity \(|\vec{v}|\) is in resonance and will thus be transferred to the excited state. If the laser beam is reflected by a mirror as in figure 5.1, on the way back the laser beam will be in resonance with atoms of opposite velocity \(-\vec{v}\), since \(\vec{k}\) has turned to \(-\vec{k}\).

If the laser beam is on resonance, both the incoming laser beam (also called saturation beam) and the reflected beam (probe beam) are in resonance with the same atoms of zero velocity. Due to that, the probe beam experiences less absorption which becomes apparent as a little dip in the Doppler-profile called Lamb-Dip (bottom of figure 5.1), which corresponds to the natural lineprofile (only power and pressure-broadened).

Figure 5.2.: Illustration of Doppler-free laser spectroscopy. On the left atoms with different velocities are addressed, which results in a strong absorption signal. On the right the same atoms with zero velocity are addressed, which results in a lower absorption signal, because the atoms are already excited by the incoming saturation beam.

Cross-over resonances

A special feature occurs for atomic systems with not just one but two excited states, whose energy difference is smaller than the Doppler-width. Performing Doppler-free saturation spectroscopy on such a system, one would expect an extended absorption-profile with 2 Lamb-dips (left of figure 5.3). Curiously, a third peak exactly in the middle of the two Lamb-Dips occurs.

The mechanism leading to such a cross-over resonance is illustrated on the right in figure 5.3.

A saturation beam with frequency \(f_l = f_1 + \frac{f_2 - f_1}{2}\) will excite those atoms to state 1, that are traveling in the same direction and have a velocity \(v\) that fulfills \(\Delta f_D = -\frac{1}{2\pi} \vec{k} \cdot \vec{v}\). The returning
probe beam is now resonant with the second transition of the same atoms. Since these atoms already absorb photons from the saturation beam, the absorption of the probe beam is reduced, which results in an additional lamb dip, referred to as cross-over resonance.

5.8. Acoustooptical modulator (AOM)

In atom-optics an AOM is a standard device for imposing a frequency shift $\Delta f$ on a given laser beam of frequency $f_{\text{in}}$. It consists of a piezo-element tapping against a crystal with the driving frequency $\Delta f$, inducing a sound wave in the crystal which acts as a diffractive grating. If the light enters with the Bragg angle

$$\sin(\theta_B) = \frac{c \cdot \Delta f}{2v_{\text{sound}} \cdot f_{\text{in}}}$$

where $c$ is the speed of light and $v_{\text{sound}}$ the speed of sound in the crystal, then part of the light is deflected. Regarding the sound wave in the crystal as phonons, that interact with the incoming photons, energy conservation requires that the frequency of the outgoing diffracted laser beam is

$$f_{\text{out}} = f_{\text{in}} + \Delta f.$$

As shown above, the angle of the deflected beam depends on the piezo-frequency $\Delta f$. This implies that if a diffracted beam is used for subsequent beam alignment, $\Delta f$ cannot be changed anymore without destroying it. This problem is overcome by an AOM used in double-pass configuration. In this configuration a retroreflector is placed behind the AOM. This way, all beams leaving the AOM are sent back to where the came from independent of the diffraction angle (i.e. independent of $\Delta f$). On their way back through the AOM they are again diffracted, such that they obtain twice the frequency shift $\Delta f$. Such a retroreflector can for example be realized by a curved mirror which is placed at the distance of its radius behind the AOM as shown in figure 5.4. Since the active region of the AOM (in this experiment 1 mm) is often smaller than the laser beam (in this experiment 2 mm), the beam is focused into the AOM by a lens. In order to obtain a low divergence of the beam at the entrance of the AOM, the focal length of the lens is usually chosen rather high (in this experiment 30 cm).
5.9. Optical fiber

An optical fiber serves as a waveguide for light. A *multi-mode* fiber consists of a core of typically 50 µm to more than 1000 µm surrounded by a cladding whose refractive index \( n_2 \) is smaller than the refractive index \( n_1 \) of the core. Rays that enter the fiber core with an angle below the *acceptance angle* are totally reflected within the core, whereas rays with a steeper angle are refracted into the surrounding cladding. Due to the large core, several modes can propagate through the fiber which leads to a wide and structured output beam profile, depending on the input (top of right figure\(^1\)).

A *single-mode* fiber has a much smaller core than a multimode fiber of a few µm, which is typically only a few times bigger than the wavelength. Light can only propagate through the fiber in a single transverse mode, which leads to a nearly Gaussian beam profile at the output.

Often a collimated laser beam has to be coupled into a fiber, which is typically done with a lens, that focuses the beam onto the core. The best coupling is achieved, if the convergence of the focused beam matches the acceptance angle and if the focused spot size matches the core-diameter of the fiber\(^2\). Therefore, the correct lens has to be chosen for each laser beam diameter. If the laser beam is too large for the lens, it will enter the fibre with an angle larger than its acceptance angle. If it is too small, the focused spot size will be bigger than the fibre core.

In this experiment a single-mode fiber is used to guarantee a perfect overlap of the probe- and couple-beam and to obtain a Gaussian beam profile.

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\(^1\)in opposite to the step-index multimode fiber shown here, a *gradient-index* multimode fiber has a gradual profile of the refractive index to avoid modal dispersion.

\(^2\)more precisely, the focused spot size has to match the *mode field diameter (MFD)* of the fibre, which is usually slightly bigger than the core.
6. Course manual

This section shall serve as a rough guideline through the experiment. Main practical information, explanation and help will be given by the tutor.

The different steps of the practical course go along with the different parts of the experiment as described in the experimental setup.

1. Spectroscopy

The aim of this part is to obtain a spectroscopy of the D1-line of Rubidium, i.e. a graph absorption vs. frequency where the absorption lines are visible. For that reason a Doppler-free saturation spectroscopy as described in section 5.7 has to be set up. The evaluation consists in

- creating a comprehensive graph of the spectrum with a calibrated frequency axis
- addressing different spectroscopy features to the corresponding transitions
- assigning values of the laser-current to the different transitions and a rough calibration laser frequency vs. laser current
- relation \( (2.16) \) between the broadened linewidth and the intensity of the laser beam shall be measured. From the resulting fit, the saturation intensity can be deduced and compared with the literature value.

2. Double-pass AOM

A double-pass AOM is needed for the couple beam as well as for the probe beam, to scan the frequency of the probe beam around the resonance frequency. The one for the couple beam is already aligned, while the one for probe has to be set up.

The AOM has a maximum diffraction efficiency of \( \sim 90\% \). Thus, for the double-pass AOM a maximum efficiency of \( \sim 80\% \) can be achieved. Efficiencies around 65\% are typically achieved in the current setup due to non-perfect laser beam properties and reflection losses at each optical device. The final efficiencies of the alignment for probe and couple should be recorded.

3. Fibre coupling

The couple beam is already coupled into the fibre, while the probe beam still needs to. With a well collimated Gaussian laser beam with a diameter that fits to the lens of the fibre coupler, more than 80\% of the laser power can be transported through the fibre. In the current setup efficiencies of only 50\% are typically achieved again due to non-perfect laser beam properties. The final efficiencies of the alignment for probe and couple should be recorded.
4. Observation of EIT

Once the previous parts have been accomplished, the main aligning work is done and it can be proceeded to the actual experiment. Probe and couple beam are leaving the fibre coupler superposed but with orthogonal linear polarization and need to be guided through the EIT cell. Behind the cell, the two beams have to be separated again with the Glan-Taylor Prism. The transmitted probe beam is imaged onto a photodiode. Optionally, the deflected couple beam can also be monitored with another photodiode. The cell needs to be heated to \( \sim 60 \degree \text{C} \) for obtaining a reasonable EIT-signal.

Some preparations have to be done:

- The \( \lambda/4 \) waveplates need to be aligned, such that the transmission of the probe beam through the Glan Taylor prism is maximal.
- The frequency of the laser has to be adjusted to the right frequency by the laser current.
- The frequency scan of the probe beam needs to be set at the AOM-driver.

If the frequencies of the AOM are set correctly and if the double-pass AOM has been adjusted to the first order, a beating signal should be seen on the photodiode due to the interference of the probe beam and the remaining contribution of the couple beam. By slightly varying the laser frequency and the power balance between probe and couple, an EIT signal (i.e. a transparency window in the transmission) should appear on the photodiode. As described in the theory section, center frequency and width of the EIT-signal depend on magnetic field and couple power, respectively. These two relations shall be measured. The evaluation should include the following:

- From relation (2.36) between center frequency and magnetic field strength Bohrs magneton and offset-magnetic field can be determined quite precisely.
- It is quite hard to determine the width of the EIT-signal with the oscilloscope. For more precise analysis the signals from the oscilloscope should be transferred to a USB-stick and evaluated with a computer using relation (2.34).
A. Appendix

A.1. Rubidium 85 D1 transition

Figure A.1.: Hyperfine structure of D1 transition in Rubidium 85, with frequency splittings between the hyperfine energy levels (taken from [21])
A.2. Rubidium 87 D1 transition

Figure A.2.: Hyperfine structure of D1 transition in Rubidium 87, with frequency splittings between the hyperfine energy levels (taken from [22]).
### Specifications

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**Note:** Performance vs Wavelength

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*Active Aperture: Aperture over which performance specifications apply.
B. Bibliography


