I. ABSTRACT

Rabi oscillations are one of the hallmarks of a quantum system. In our experiment we measure Rabi oscillations at room temperature between the $m_s = 0$ and $m_s = 1$ ground state spin levels of the nitrogen-vacancy center. We will also describe a double-pass acousto-optic modulator setup which we built to operate as a fast light shutter, a crucial part of the system to be used for the Ramsey sequence. We will use these measurements to determine the width of the $\pi/2$ pulse which is needed in order to perform Ramsey and spin-echo sequences.

II. MAGNETOMETRY OF NITROGEN VACANCY CENTERS

A. Introduction

The microscopic model and most optical properties of ensembles of nitrogen-vacancy (NV) centers were firmly established in the 1970s based on the optical measurements combined with structural distortion and on electron paramagnetic resonance (EPR) [1]. However, a minor error in EPR results (it was assumed that illumination is required to observe NV EPR signals) led to incorrect multiplicity assignments in the energy level structure [2]. In 1991 it was shown that EPR can be observed without illumination [3]. The characterization of single NV centers has become a very competitive field nowadays, with many dozens of papers published in the most prestigious scientific journals. One of the first results was reported in 1997 [4]. In that paper, it was demonstrated that the fluorescence of single NV centers can be detected by room-temperature fluorescence microscopy and that the defect shows strong photostability. Also one of the outstanding properties of the NV center was demonstrated, namely room-temperature optically detected magnetic resonance.
The detection of single centers soon enabled demonstrations of photostable single photon generation, which highlighted the NV center for implementation in quantum optical networks [5], as well as demonstrations of optical preparation and readout of the electronic spin of the centers, which identified them as a possible solid-state spin qubit suitable for quantum information processing and quantum metrology devices. Following these studies, the growth of research into the NV center and the development of applications have been incredibly rapid and a number of important milestones have been reached. Beyond quantum information processing, there exist important applications in room-temperature nanoscale magnetometry, bio-magnetometry, electrometry, decoherence microscopy and, recently, using the NV center as a sensitive nanoscale thermometer [6].

B. Crystalline Structure and Production of NV centers

The negatively charged NV center is formed by a missing carbon adjacent to a nitrogen ion impurity in the diamond lattice (see Fig. 1). This ionic structure is accompanied by a spatial electronic distribution that can be represented as a single wavefunction isolated from the crystal electronic wavefunction (responsible for the diamond band structure) [7]

![Spatial configuration of an NV center in a diamond unit cell](image)

FIG. 1: Spatial configuration of an NV center in a diamond unit cell, where the $C_{3v}$ symmetry is easily observed. Carbon atoms are represented by gray spheres, nitrogen in blue and vacancy in red.

Nitrogen vacancy centers are typically produced from a single nitrogen atom impurity in place of a carbon atom (called C or P1 centers in diamond literature) by irradiation followed by annealing at temperatures above 700°C. A wide range of high-energy particles are suitable for such irradiation, including electrons, protons, neutrons, ions, and gamma-ray photons. Irradiation produces lattice vacancies, which are a part of the NV centers. Those vacancies are immobile at room temperature, and annealing is required to move them. Single substitutional nitrogen produces strain in the diamond lattice. It therefore
efficiently captures moving vacancies, producing the NV centers. During chemical vapor deposition of diamond, a small fraction of single substitutional nitrogen impurity (typically < 0.5%) traps vacancies generated as a result of the plasma synthesis. Such nitrogen vacancy centers are preferentially aligned in the growth direction. Diamonds are notorious by having relatively large lattice strain. Strain splits and shifts optical transitions from individual centers resulting in broad lines in an ensemble of centers [1]. Special care is taken to produce extremely sharp NV lines (line width ≈ 10 MHz) required for most experiments: high-quality, pure natural or better synthetic diamonds (type IIa) are selected. Many of them already have sufficient concentrations of grown-in NV centers and are suitable for applications. If not, they are irradiated by high-energy particles and annealed. Specific irradiation dosages are chosen to tune the concentration of NV centers produced such that individual NV centers are separated by micrometer order distances. Individual NV centers can then be studied with standard optical microscopes or with near-field scanning optical microscopes that can provide sub-micrometer resolution.

C. Energy Level Structure and Optically Detected Magnetic Resonance (ODMR)

Sensing of magnetic and electric fields is based on optically detected resonances, utilizing laser light and micro-waves (MW) [8–10]. The NV center has a spin triplet ground state (g.s.) with a zero-field splitting of 2.87 GHz. The \( m_s = \pm 1 \) g.s. energy levels are degenerate in the absence of a magnetic field. When this degeneracy is lifted (Zeeman splitting) a detectable change in the NV center spectroscopy provides the ability to measure the magnetic field. The NV center energy-level structure is depicted in Fig. 2.

![Nitrogen vacancy center energy-level structure. Radiative transitions are indicated by solid arrows and non-radiative transitions by dashed arrows.](image)

The g.s. spin can be determined by laser excitation followed by measuring the fluorescence...
cence of the spontaneous emission. If the spin projection on the NV axis is zero, the NV center will be excited to the $m_s = 0$ excited state and will then relax to the $m_s = 0$ g.s., emitting a 637 nm photon. If the NV center is in the $m_s = \pm 1$ state it will be excited to an $m_s = \pm 1$ excited state and will then have a $\sim 20\%$ probability to relax to the $m_s = 0$ g.s. via non-radiative decay. This enables to optical pumping of the NV’s to the $m_s = 0$ g.s. and measuring the NV’s spin state by measuring its fluorescence. After pumping the NV’s to the $m_s = 0$ g.s., it is possible to excite them to the $m_s = \pm 1$ g.s. with MW radiation. The MW resonance frequency is proportional to the external magnetic field (the energy levels are Zeeman shifted). An example of the NV spectrum (measured in our lab) with and without a magnetic field is depicted in Fig. 3.

![ODMR of a [1,1,1] diamond without external magnetic field](image)

![ODMR of a [1,1,1] diamond with external magnetic field](image)

**FIG. 3:** Nitrogen vacancy spectroscopy measured in our lab without (upper) and with (lower) an external magnetic field. The magnetic field is applied in a general (not along any specific axis) direction with a $\sim 64$ gauss component along the [1,1,1] crystallographic axis, clearly exposing all existing orientations.

### III. BLOCH SPHERE REPRESENTATION

The state of an atomic two-state system (and any other two-state system) can be represented as a superposition of the two eigenstates of the Hamiltonian $H^0$ describing the
system:
\[ |\Psi\rangle = C_0 |0\rangle + C_1 |1\rangle, \]
where \( C_0 \) and \( C_1 \) are complex numbers such that \( \sum_{i=0}^{1} C_i C_i^* = 1 \) and additionally that
\[ H^0 |0\rangle = E_0 |0\rangle, H^0 |1\rangle = E_1 |1\rangle. \]

We can also write Eq.(1) as follows:
\[ |\Psi\rangle = \cos \left( \frac{\theta}{2} \right) |0\rangle + e^{i\phi} \sin \left( \frac{\theta}{2} \right) |1\rangle \]
where \( \theta \) and \( \phi \) are the spherical coordinates for the bloch unit-vector:
\[ \vec{a} = (\sin \theta \cos \phi, \sin \theta \sin \phi, \cos \phi). \]

As long as our system does not interact coherently with any external system, a global phase can be neglected in Eq.(3).

A useful geometrical interpretation of a two-state system is the “Bloch Sphere”, a three-dimensional unit sphere on which quantum states are represented as points (Fig. 4). The ”north pole” of the sphere (\( \theta = 0 \)) denotes the pure state \( |0\rangle \) and the ”south pole” (\( \theta = \pi \)) denotes the pure state \( |1\rangle \). All other points represent superposition states according to Eq. (1).
IV. RABI OSCILLATIONS

The population of a two-state system initially prepared in the lower state and then subjected to resonant (or slightly detuned) radiation, will undergo oscillations between the two states, usually referred to as Rabi oscillations. The resonant Rabi frequency $\Omega_0$ is proportional to the intensity of the incident radiation, and as we detune the radiation from resonance, the frequency of the oscillations increases while their amplitude decreases. The decay of the Rabi oscillations over time is characterized by the ensemble coherence time $T_2^*$. Therefore, tracking the Rabi oscillations of our system should provide data regarding its coherence time as well as the radiation power felt by the NV centers. Hence, we probed the Rabi oscillations in our system between the $m_s = 0$ and one of the $m_s = \pm 1$ levels, the latter being separated by an external magnetic field. In the experiment (see Fig. 5), the diamond is illuminated constantly by green light at low laser power, which has a weak pumping effect but which still induces fluorescence proportional to the spin state population. Peak fluorescence corresponds to concentrating all the population in $m_s = 0$, while minimum levels of fluorescence correspond to full population transfer to $m_s = \pm 1$. We simultaneously apply short microwave pulses ($\approx 20 \mu s$) (Fig. 6). As a result, the system is pumped (pumping

FIG. 5: Schematic layout of the experiment

FIG. 6: Pulse sequence for our Rabi Oscillations
time was $\sim 30 \text{ ms}$) to the $m_s = 0$ level when the microwave radiation is turned off and Rabi oscillations occur when the MW is on. Synchronizing the oscilloscope with the microwave pulse displays the evolution of the oscillations on our screen. In Fig. 7 we display Rabi oscillations between $m_s = 0$ and $m_s = -1$ measured in the lab. We fit the data to the function

$$F(t) = A [1 - B \cos (\Omega t + \phi)] e^{-\frac{t}{\tau}} + C$$

where $F(t)$ is the fluorescence as a function of time, and $\tau, \Omega$ are the ensemble coherence time and the frequency of the oscillations, respectively.

![FIG. 7: Rabi oscillations measured in the lab](image)

From the curve fit we deduced:

$$T_s^* = 0.21 \pm 0.02 \ \mu s$$

$$\Omega = 2\pi \times 1.19 \pm 0.01 \text{ MHz}$$

V. DOUBLE-PASS ACOUSTO-OPTIC MODULATOR (AOM)

A. Introduction

Acousto-optic modulators (AOM’s) are widely used to accomplish frequency control in laser cooling experiments. When the laser frequency is scanned with an AOM, the angle of the first order diffracted beam also shifts, since the beam diffraction angle is a function of the modulation frequency. Changes in the beam direction may be desirable for some applications where spatially resolved diffracted beams are needed, but for many applications any change in the laser propagation direction is an unwanted side effect. Using an AOM in the double-pass configuration is a way to eliminate changes in beam steering during frequency sweeps and jumps within the frequency tuning bandwidth of the AOM. We will briefly discuss the theory and how to practically construct a system utilizing a double-pass AOM.
B. Bragg scattering description

For most practical applications of AOM’s, the Bragg description of the modulation process is a good approximation to the behavior of the system. The main features of AOMs are derived in this picture by treating the modulation as a photon-phonon scattering process. The Bragg matching condition can be derived by treating the acoustic and optical fields as particles with momentum $\kappa$ and $k$, respectively, where $\kappa(k)$ is the phonon (photon) wave vector for the acoustic (optical) field, given by $\kappa=\Omega/v_s$ where $\Omega$ is the radio-frequency (RF) modulation frequency and $v_s$ is the speed of sound in the crystal. Similarly, $k=\omega/v_L$, where $\omega$ is the light frequency and $v_L$ is the speed of light in the crystal. A scattering process between photons and phonons results in the absorption or emission of acoustic phonons. A first-order scattering process between a photon and a single phonon is described by the energy momentum relations:

$$\omega_d = \omega_i \pm \Omega$$

$$k_d = d_i \pm \kappa$$

The subscripts $i$ and $d$ designate whether the corresponding photon is incident or diffracted. The sign denotes phonon absorption or emission, depending on the relative orientations of the incident photon and phonon wave vectors. The Bragg matching condition that determines the optimum angles for the incident laser and acoustic beams for peak first-order diffraction efficiency can be derived with energy and momentum conservation arguments. Fig. 8 shows momentum conservation diagrams which describe photon-phonon scattering events in which a phonon is absorbed (+1st-order diffraction). The photon momentum of the diffracted light is equal to the sum of the momenta of the phonon and the incident photon. Conservation of energy requires that the frequency of the diffracted beam be shifted upward from $\omega$ to $\omega + \Omega$ for a phonon absorption process; but since $\omega \gg \Omega$, the frequency shift can be ignored in the momentum conservation analysis, and $|k_i| = |k_d|$. Adding together the $x$ and $y$ momentum components leads one to Bragg’s Law- $\sin \theta_B = \kappa/2k_i$ where $\theta_B$ is the Bragg angle and $\theta_i = \theta_d = \theta_B$.

Note that we have derived this equation without considering the effect of the boundaries of the acoustic medium. Often it is desirable to calculate the Bragg angle outside of the crystal, $\theta_{B,\text{Ext}}$. To a close approximation, if the crystal boundaries are parallel to $\kappa$, the external Bragg angle is larger than the internal Bragg angle by a factor of $n$, where $n$ is the refractive index of the acoustic medium. The cases shown in Fig. 8 represent the two
FIG. 8: Momentum conservation diagrams for the absorption of a phonon. The two cases [(a) and (b)] represent two configurations for the incident and diffracted light beams where the conditions are Bragg matched for a given direction of the phonon momentum. Note that $\kappa_i$ and thus $\theta_i$ and $\theta_d$, are exaggerated in the diagrams for clarity. The Bragg angle is normally about $1^\circ$ or less, and the phonon momentum is roughly two orders of magnitude smaller than the photon momentum [11].

geometries where the Bragg matching condition is met for +1st-order phonon absorption for a given direction of $\kappa$. Notice that the incident photon wave vector for one case and the diffracted photon wave vector for the other case counter-propagate. This point is relevant to the discussion of the double-pass configuration below. Diagrams similar to those in Fig. 8 can be drawn for the single phonon emission process, where the diffracted light enters the $-1$st-order, and the frequency of the diffracted beam is shifted to $\omega - \Omega$.

C. **Acousto-Optic Devices**

A schematic drawing of an AOM is shown in Fig. 9. An RF signal is fed to a strain transducer in contact with the AOM crystal. The RF modulation at frequency $\Omega$ causes a traveling density wave to form inside the crystal. the wave propagates at the speed of sound in the crystal, $v_s$ with the frequency $\Omega$. The refractive index is therefore modulated with a wavelength of $\Lambda = 2\pi v_s/\Omega$, and the crystal acts like a thick diffraction grating with the rulings traveling away from the transducer with a velocity $v_s$. The Bragg approximation is valid only when the acoustic wave may be described as a plane wave and all phonons have the same wave vector. In practice, this limiting case can be achieved to a close approximation when the strain transducer is long compared to the acoustic wavelength in the direction of laser beam propagation and acoustic diffraction is minimized. It is therefore not uncommon to achieve greater than 80% diffraction efficiency into a single diffraction order with careful engineering.
FIG. 9: Schematic drawing of an aligned AOM. The AOM is shown in the mode where the light is diffracted into the first order and the laser frequency is up-shifted by the RF modulation frequency Ω. Note that phonons are absorbed in this configuration, and momentum is transferred to the laser beam in the direction away from the strain transducer. The crystal can also be aligned to be Bragg matched for the $-1^{st}$-order (not shown). In that case, corresponding to a phonon emission process, all of the laser beams are mirror symmetric about the horizontal axis with respect to those shown in the figure, and the diffracted laser frequency is down shifted by Ω. Note that the external Bragg angle is larger than the internal Bragg angle because of refraction at the crystal surfaces [11].

D. Double-Pass Configuration

Using a double-pass AOM setup, in which the laser beam goes through the AOM twice, allows us to significantly lower the power of the beam very quickly, thus creating an on/off switch for the laser. The configuration we used is described in a schematic drawing in Fig. 10 and also photographed in Fig. 11. Modulating the light with the AOM twice also means that the total frequency shift is twice the modulation frequency.

FIG. 10: Schematic drawing of a double-pass AOM system

The two momentum conservation diagrams in Fig. 8 represent the momentum vectors for the two passes. Notice that the incident and diffracted beams for one case counter-propagate the diffracted and incident beams for the other case. Thus with the alignment optimized, if the $+1^{st}$-order diffracted beam is reflected back onto itself through the AOM for a second pass, the $+1^{st}$-order diffracted beam for the second pass will spatially overlap the original incident beam, independent of the modulation frequency, and the direction of the laser beam can thereby be made independent of its frequency shift. In our experiment, the incident beam is filtered out using a polarizing beam splitter and a quarter wave plate (Fig. 11) and so only the beam which undergoes a double pass continues towards the diamond.
FIG. 11: Photo of our double-pass AOM setup with added labels and beam trace [QWP: quarter-wave Plate, PBS: polarizing beam splitter]
VI. REFERENCES


