A quantum gas with tunable interactions in an optical lattice

Dissertation by Mattias Gustavsson



submitted to the Faculty of Mathematics, Computer Science and Physics of the University of Innsbruck

> in partial fulfillment of the requirements for the degree of doctor of science

Advisor: a. Univ. Prof. Dr. Hanns-Christoph Nägerl

Innsbruck, December 2008

Abstract

The combination of an ultracold gas with a periodic potential in the form of an optical lattice opens up the opportunity to study phenomena known from solid state physics in a clean and well isolated environment, with a high degree of control over both internal and external degrees of freedom. This thesis reports on the realization of a tunable quantum gas in an optical lattice, where the use of Cs atoms allows a precise control over the atomatom interactions using a broad Feshbach resonance. In particular, it is possible to strongly suppress the interactions by tuning the scattering length close to zero.

In the framework of this thesis, a Cs BEC apparatus was constructed with the specific aim to perform experiments with optical lattices. The apparatus is an evolution of the first-generation Innsbruck Cs BEC apparatus. Instead of using a stainless steel vacuum chamber, the Cs atoms are trapped in a glass cell, which allows for fast and precise control over magnetic fields without disturbing eddy currents. The setup was designed to allow a large optical access, enabling the addition of an optical lattice, and is capable of producing a Cs BEC of up to $2 \cdot 10^5$ atoms every 10 s.

The control over atom-atom interactions is demonstrated in two sets of experiments studying the effect of interactions on a Bloch oscillating BEC. The atom-atom interactions lead to density-dependent phase shifts at the individual lattice sites and limit the number of Bloch oscillations one can observe. In the first set of experiments, we quantitatively characterize this dephasing as a function of the magnetic field and determine the point where atom-atom interactions are minimized. With interactions minimized, more than 20000 Bloch oscillations can be followed, corresponding to a coherent evolution over more than 10 s. The force inducing the Bloch oscillations can then be determined with better than 10^{-6} precision. Our technique to suppress interactions has potential applications for BEC atom interferometry, where phase shifts and decoherence due to interactions are a major problem.

In the second set of experiments, we observe and control matter wave interference that is driven by interparticle interactions. We show that interaction-induced phase shifts lead to the development of a regular interference pattern in the wave function of a Bloch oscillating BEC. The high degree of coherence in this process is demonstrated in a matter wave spinecho type experiment, where the phase evolution of a dephased BEC is reversed by tuning the scattering length close to zero and applying an external potential, allowing us to recover the original BEC wave function.

Acknowledgments

It is often said that no matter how interesting a thesis is, the acknowledgments is the part that is read first. This is as it should be. In our branch of experimental physics, work is done in teams and the work presented in this thesis is certainly no exception. It would have been neither possible, nor enjoyable, had it been done alone. I would here like to acknowledge the contributions of a number of colleagues, friends and family.

Hanns-Christoph, my advisor, gave me the opportunity to write this thesis. His warm and caring personality, enthusiasm and skills made it a joy to work on the project, and he is always available and listens when something comes up. In addition, he gives you a lot of freedom to pursue your own ideas, something I very much appreciate.

Rudi, our group leader, was bold enough to invite an unknown guy from far away Sweden with very little experimental experience and give him the chance to prove himself in the lab. Rudi has created a world-class research environment and works very hard to give us the possibility to do great science.

Elmar shared the neverending fight with technology and laws of nature for the sake of peering a little further into the unknown. His steady and methodical character was the perfect addition to our team, and I am also thankful for his (mostly) very good taste in choosing lab music. We shared both good and bad times in the lab and I value him a lot, both as colleague and as friend. May you never have to touch an air conditioning unit again.

Apart from Elmar, there have been many other members of the Cs III team over the years, all contributing to a fun and inspiring lab environment and also sharing good times outside the lab. **Peter** set up the diode lasers and taught me mountain biking. **Toni** helped winding coils and setting up the laser cooling. **Gabriel** did a very careful job with the optical lattice. **Manfred** built lasers for the future BEC interferometer and applied his electronic wizard skills to all parts of the experiment. **Hans** is a living molecule encyclopedia and leads the effort to create deeply bound molecules. **Russ** brought in knowledge from the other side of the Atlantic. I leave knowing that the apparatus is in very good hands. Thanks guys, it's been a pleasure working with you.

Stefan R has been a steady mountain bike companion, patiently waiting for me downhill so I could return the favour uphill. He has an uncanny sense for finding great trails ending in swamps, creeks or other impassable terrain. Thanks for all the good chats!

The Grimm group is a wonderful and diverse bunch of people. We shared ideas, equipment and free time, be it at parties, barbecues, skiing, sledding or biking. Additionally, we proved every year in the Stadtlauf that experimentalists run much faster than theorists.

The supporting administrative and technical staff are essential to keep things running smoothly. **Christine, Karin, Patricia, Gabriel, Nicole, Manuel, Arthur, Gerhard, Toni, Josef, Helmut** – thanks for making life easier in so many ways.

Many friends made life in Innsbruck enjoyable and provided support when it was not. Among them were **Petra**, **Régis**, **Andrew D**, **Olof**, **Andrew M**, **Anna**, **Valeria** and all the members of my beloved Tyrolean marching band, **Musikkapelle Allerheiligen**.

My family is my biggest support. Mum, dad and Emelie, thank you for always being there.

Again, many thanks to all of you. It's been a great voyage.

Contents

1	Intro	troduction						
2	A new Cs BEC apparatus							
	2.1	Basic concepts	13					
		2.1.1 Ultracold collisions	13					
		2.1.2 Bose-Einstein condensation	16					
	2.2	Experimental setup	20					
		2.2.1 The Cs atom	20					
		2.2.2 Vacuum system	21					
		2.2.3 Magnetic fields	26					
		2.2.4 Zeeman slower	30					
		2.2.5 Magnetic levitation	31					
		2.2.6 Diode laser system	32					
		2.2.7 Experiment control	33					
		2.2.8 Detection and diagnosis	35					
	2.3	The path to BEC	37					
		2.3.1 Magneto-optical trap	38					
		2.3.2 3D Raman sideband cooling	39					
		2.3.3 Large volume dipole trap - reservoir	41					
		2.3.4 Tightly focused dipole trap - dimple	43					
		2.3.5 Evaporation and Bose-Einstein condensation	44					
3	A BEC in an optical lattice 4							
	3.1	Theory	47					
		3.1.1 Optical lattice potential	47					
		3.1.2 Bloch states and band structure	48					
		3.1.3 Wannier states	50					
		3.1.4 Effective 1D equation	53					
		3.1.5 Ground state of a BEC in a lattice	55					
	3.2	Technical setup	57					
		3.2.1 Lattice setup	57					
		3.2.2 Lattice depth calibration	59					
4	Bloo	h oscillations and interaction-induced dephasing	63					
	4.1	Theory of Bloch oscillations	64					
	4.2	Experimental realization	65					

Contents

	4.3	4.2.1 4.2.2 4.2.3 4.2.4 Possib 4.3.1 4.3.2	Observation of Bloch oscillations	65 67 68 71 73 73 76				
5 Coherent dephasing of Bloch oscillations								
	5.1	Theory	y	81				
		5.1.1	Dephasing in the limit of a strong force	81				
		5.1.2	Shift in oscillation period due to the lattice position	83				
		5.1.3	From the 1D model to a 2D absorption picture	84				
	5.2	Experi	mental realization	84				
		5.2.1	Experimental parameters	84				
		5.2.2	Structure of a dephased cloud in quasimomentum space	85				
		5.2.3	Cancelation of dephasing through an external potential	86				
		5.2.4	Rephasing of a dephased condensate	92				
		5.2.5	Decay and revival of Bloch oscillations	92				
6	Out	ook		97				
Bi	Bibliography							
A Publications								

1 Introduction

The invention of laser-cooling techniques gave birth to the field of ultracold gases, where dilute samples of atoms can be prepared at microkelvin temperatures. Evaporative cooling made it possible to achieve even lower temperatures and reach the quantum gas regime. Here, the motion of the atoms cannot be described by classical physics and a quantum mechanical description has to be used, with the atoms treated as wavepackets. For bosonic atoms, a phase transition occurs when the temperature is so low that the extent of the wavepackets, characterized by the de Broglie wavelength, becomes comparable to their average mutual distance. A macroscopic number of particles accumulate in the quantum mechanical ground state of the trapping potential, a phenomenon first predicted by Bose and Einstein in the early days of quantum mechanics [Bos24, Ein25]. The experimental realization of such a Bose-Einstein condensate (BEC) in a dilute gas of ⁸⁷Rb [And95] in 1995 was a major breakthrough that was rewarded with the Nobel prize six years later. In the years that followed, the notion of a BEC as a coherent matter wave was validated via interference of two independent condensates [And97], and superfluidity was proven through the excitation of vortices [Mat99, Mad00, AS01]. The nonlinear behaviour of the BEC matter wave was demonstrated with the realization of a matter wave amplifier [Ino99] and the excitation of solitons [Den00].

By now, Bose-Einstein condensation has been reached with all alkali atom species except Fr, and additionally with H, Yb, Cr, metastable He and with weakly bound molecules of Li_2 and K_2 . The different species all have different advantages. For example, the large internal energy of metastable He allows single atom detection, whereas Cr has a large magnetic dipole moment, allowing the investigation of dipolar effects in quantum gases.

The heaviest alkali atom, Cs, offers very attractive scattering properties. The presence of both broad and narrow Feshbach resonances at moderate magnetic fields [Chi04] enables a very precise control over the atom-atom interactions, characterized by the s-wave scattering length. It therefore allows the creation of a tunable quantum gas, where the interactions can be precisely varied from repulsive to attractive, and even be set very close to zero [Web03b]. Feshbach resonances also make it possible to create an ultracold molecular sample out of a trapped atomic gas [Reg03, Her03].

The combination of an ultracold gas with a periodic potential in the form of an optical lattice opens up new exciting opportunities. Many of the phenomena pertinent to solid state

1. Introduction

physics can be investigated in a clean system well isolated from the environment, with a lattice potential free of defects. Light fields, radio frequency and microwave radiation and magnetic fields can be employed to control both the motion and the internal state of the atoms, giving a large degree of control over different model parameters. Absorption imaging allows the experimentalist to directly measure density and momentum distributions, and it is also possible to obtain information on spatial correlations. The realization that bosonic atoms in an optical lattice can be described by a Bose-Hubbard model [Jak98] and the subsequent observation of the quantum phase transition from a BEC to a Mott insulator [Gre02] launched the field of strongly interacting quantum systems, and the Mott insulator state with a well-defined number of atoms per lattice site is an ideal laboratory for few-body physics.

In this thesis, I report on the realization of a *tunable* quantum gas in an optical lattice. In contrast to previous experiments with optical lattices, our use of Cs atoms offers very precise control over the interparticle interactions and therefore further control over the system parameters than previously possible. This addition of a new experimental "knob" to turn allows us to investigate the role of interactions in different physical systems, both in the weakly and the strongly interacting regime. In particular, we study the effect of interactions on a Bloch oscillating BEC in two series of experiments. The well-known phenomenon of Bloch oscillations [Blo29, BD96, And98] originates when a wave-packet in a periodic potential is subject to a constant force, which causes an oscillatory motion. The number of Bloch oscillations one can observe is however strongly limited by collisional dephasing. In the first set of experiments, we measure the dependence of this dephasing on the interaction strength. The interparticle interactions can can be set very close to zero by tuning the magnetic field near a zero-crossing of the s-wave scattering length. Control over the scattering length to an unprecedented precision of 0.1 Bohr radii is demonstrated and the magnetic field for which interactions are minimized is determined to a high precision. With the atom-atom interactions minimized, more than 20000 Bloch oscillations can be followed.

The strong suppression of atom-atom interactions has an additional interest for atom interferometry. A BEC combines high brightness with narrow spatial and momentum spread and would constitute an ideal source for a matter wave interferometer. Unfortunately, because of the high density in a BEC, interactions lead to phase diffusion and can cause systematic frequency shifts due to unwanted density gradients, limiting the performance of the interferometer. This limitation could be overcome around using the precise control over the interaction strength demonstrated in this work.

In a second series of experiments, we observe matter-wave interference that is driven by interparticle interactions. When a large force is applied to a BEC to induce Bloch oscillations, tunneling between the lattice sites is strongly suppressed. The system can then be seen as a matter wave interferometerm, where every lattice site experiences a different phase shift proportional to the local potential at the site. As is well known, the potential gradient due to the applied force leads to Bloch oscillations. In this work, we demonstrate that the additional interaction potential leads to an additional component in the evolution of the phase at the individual lattice sites, which can be detected and visualized by the appearance of interference patterns when the wavefunction is imaged in momentum space. The high degree of coherence in the system is demonstrated by reversing the wave function evolution of a dephased BEC, switching off interactions and applying an external potential. The original BEC wave function is then recovered.

Overview

This thesis is structured in the following way. In chapter 2, the basic concepts needed to understand ultracold gases and Bose-Einstein condensates are introduced. The technical details of our new generation Cs BEC apparatus are then described, and the different cooling steps employed to produce a Cs BEC are outlined. The third chapter discusses a BEC in an optical lattice. I review the important concept of band structure and Bloch states, and also introduce the Wannier functions, an alternative basis of wavefunctions localized to single lattice sites. The Wannier functions are then used to derive an effective 1D equation describing the dynamics of a BEC in a lattice. Using these tools, the properties of the BEC ground state in the lattice are calculated. The technical setup of the lattice is then described together with an important method for measuring the lattice depth.

Chapter 4 presents the measurements of the interaction-induced dephasing of Bloch oscillations, demonstrating the very precise control over atom-atom interactions possible in our setup, particularly the ability to minimize atom-atom interactions. This opens up new possibilities for BEC interferometry, and this chapter concludes with a study arguing that a precision measurement of the fine structure constant with a BEC contrast interferometer would be feasible. In chapter 5, the evolution of the BEC wave function is studied. A simple analytical model is derived that explains the interference fringes appearing in the momentum wave function of a dephased Bloch oscillating BEC as a consequence of interaction-induced phase shifts. The appearance of these fringes and the coherence of the process is demonstrated experimentally. Finally, chapter 6 discusses some of the many exciting prospects for future experiments with a tunable quantum gas in an optical lattice.

Publications

The following articles have been written in the framework of this thesis. They are attached as part of the appendix.

Control of Interaction-Induced Dephasing of Bloch Oscillations M. Gustavsson, E. Haller, M. J. Mark, J. G. Danzl, G. Rojas-Kopeinig, and H.-C. Nägerl Phys. Rev. Lett. **100**, 080404 (2008).

Quantum Gas of Deeply Bound Ground State Molecules J. G. Danzl, E. Haller, M. Gustavsson, M. J. Mark, R. Hart, N. Bouloufa, O. Dulieu, H. Ritsch, and H.-C. Nägerl Science **321**, 1062 (2008), published online 10 July 2008; 10.1126/science.1159909.

Dark resonances for ground state transfer of molecular quantum gases M. J. Mark, J. G. Danzl, E. Haller, M. Gustavsson, N. Bouloufa, O. Dulieu, H. Salami, T. Bergeman, H. Ritsch, R. Hart, and H.-C. Nägerl Submitted for publication. arXiv:0811.0695.

Precision molecular spectroscopy for ground state transfer of molecular quantum gases J. G. Danzl, M. J. Mark, E. Haller, M. Gustavsson, N. Bouloufa, O. Dulieu, H. Ritsch, R. Hart, and H.-C. Nägerl Accepted for publication in Faraday Discuss. arXiv:0811.2374.

1. Introduction

Interference of interacting matter waves M. Gustavsson, E. Haller, M. J. Mark, J. G. Danzl, R. Hart, A. Daley, and H.-C. Nägerl Submitted for publication.

2

A new Cs BEC apparatus

To perform the experiments described in this thesis, a new apparatus for trapping and cooling Cs atoms to quantum degeneracy was developed. The goal was to set up a machine that can rapidly produce a Cs BEC which can be loaded into an optical lattice, with rapid and precise control over magnetic fields and large optical access to allow maximum flexibility for future experiments. In this chapter, I will first review the basic concepts needed to understand cold trapped atoms and Bose-Einstein condensates. The technical setup of the apparatus will then be described, and finally the different cooling steps necessary to achieve BEC will be detailed.

2.1 Basic concepts

2.1.1 Ultracold collisions

As is well-known, scattering can be treated by expanding the wave function of the relative motion of two colliding atoms in spherical partial waves. Each of the spherical partial waves is characterized by its angular momentum, *l*. At sufficiently low energies, the centrifugal barrier prohibits partial waves with nonzero angular momentum and only *s*-wave scattering (l = 0) needs to be considered. The scattering is then isotropic and is characterized by the phase shift δ_0 between the incoming and the outgoing *s*-wave. In the limit of zero collision energy, the scattering is usually parameterized by the scattering length

$$a = \lim_{k \to 0} \frac{\tan \delta_0(k)}{k},\tag{2.1}$$

where *k* denotes the wave vector of the relative motion of the atoms. The scattering behaviour in the low-energy limit is thus well described by one single parameter, the scattering length, independent of the details of the interaction potential between the two colliding particles. The s-wave scattering length is typically in the range $10 - 100 a_0$ for alkali atoms, where a_0 is the Bohr radius.

For two identical bosons, the collisional cross-section is [Dal99b]

$$\sigma_{\rm el} = \frac{8\pi a^2}{1 + k^2 a^2}.$$
(2.2)

This expression has two limiting cases. For large scattering lengths, such that $ka \gg 1$, the cross section is limited by the collision energy, $\sigma_{\rm el} = 8\pi/k^2$. This is called the unitarity limit. In the limit of small scattering length, $ka \ll 1$, the cross section is $\sigma_{\rm el} = 8\pi a^2$.

In the s-wave limit, the exact interaction potential can be approximated by a point-like scattering potential. This contact interaction potential reads ¹

$$V(r) = g\delta(r), \tag{2.3}$$

where r is the distance between between the colling particles and the coupling constant g is proportional to the scattering length,

$$g = \frac{4\pi\hbar^2 a}{m}.$$
(2.4)

Besides the contact interaction, there can also be other more long-range interactions where the scattering potential cannot be described by a δ -function, for example different forms of dipole-dipole interaction. In most cases the contact interaction dwarfs the other forms of interaction, but there are examples where this is not the case. Chromium has a large magnetic dipole moment and the magnetic dipole-dipole interaction can be made much larger than the contact interaction when the scattering length is changed using a Feshbach resonance [Lah07]. Samples of ultracold polar molecules in the rovibronic ground state, which have a very large electric dipole moment, have been realized in recent experiments [Sag05, Ni08]. It will be demonstrated in section 4.2.3 that the magnetic dipole-dipole interaction is not negligible for the atom we use, Cs, when the scattering length is tuned close to zero.

Feshbach resonances

The scattering length does not have to be constant. It can in many cases be tuned by an external magnetic field through so-called Feshbach resonances, a concept first studied in the context of nuclear physics [Fes58] and later applied to atom-atom scattering [Tie93].

The principle behind a Feshbach resonance is illustrated in figure 2.1. In the preceding discussion, we did not take into account the internal structure of the colliding particles. However, atoms do have an internal structure and the interaction potential between two particles, usually called a scattering channel, depends on their internal state. The channel corresponding to the initial state of the colliding particles is called the *incident channel*. During a collision, the atoms can change their internal state and exit in another channel, provided there is a coupling between the channels. This can only happen if the continuum of this outgoing channel has a lower energy than the total energy of the incident channel, in which case the channel is called an *open channel*. It is not possible to scatter into a channel where the continuum has a higher energy than the incident channel, and such a channel is therefore called a *closed channel*. Coupling to such a channel can however still modify the scattering properties. If the internal atomic states corresponding to the closed channel have a different magnetic moment than those of the incident channel, the "position" of the closed channel can be changed by tuning the magnetic field. A bound molecular state in the closed channel can then be brought into degeneracy with the incident scattering state. If there is some

¹Note that this simple form of the contact interaction potential is only valid when using the Born approximation for scattering. A more proper way to express the contact interaction is the potential $V(\mathbf{r})\Psi(\mathbf{r}) = g\delta(\mathbf{r})\frac{\partial}{\partial \mathbf{r}}(\mathbf{r}\Psi(\mathbf{r}))$. See [Dal99b] for more detail.



Figure 2.1: Schematic illustration of a Feshbach resonance. **Left.** Molecular potentials: An external magnetic field can be used to tune a bound molecular state in degeneracy with the scattering state of two free atoms. **Right.** Zeeman diagram: When the molecular state and the scattering state (here with zero magnetic moment) are tuned into degeneracy, the scattering length diverges.

coupling between the channels, this leads to resonant scattering and a divergence of the scattering length.

A Feshbach resonance can be characterized by its position B_0 , the magnetic field where the molecular state crosses the incident scattering state², and its width ΔB , which is dependent on the magnetic moment of the bound state and the strength of the coupling between the two scattering channels. The scattering length around a Feshbach resonance can be written as

$$a(B) = a_{\text{bg}} \left(1 - \frac{\Delta B}{B - B_0} \right), \tag{2.5}$$

where a_{bg} is the background scattering length far from resonance.

Coupling between the scattering channels can be induced by several forms of interaction. The Coulomb interaction preserves orbital angular momentum l, which for *s*-wave scattering means that only molecular *s*-wave states couple to the incident channel. These resonances are consequently called *s*-wave resonances. Relativistic interactions such as magnetic spin-spin interaction and second order spin-spin interaction are usually much weaker. They couple the *s*-wave scattering state to molecular states with higher orbital angular momentum l = 2, 4, ... and therefore give rise to *d*-wave (l = 2) and *g*-wave (l = 4) resonances.

The atom we use in our experiments, Cs, has an unusually large second-order spin-spin interaction. This leads to a very rich variety of Feshbach resonances and makes it a very interesting atom for use in cold atom experiments. The Cs scattering properties have been extensively investigated in a series of experiments at Stanford and an accompanying theoretical analysis at NIST [Chi00, Leo00, Chi04]. The scattering length for Cs atoms in the $|F = 3, m_F = 3\rangle$ state is shown in figure 2.2. For magnetic fields between 0 and 150 G, field strengths that are easily accessible in the lab, seven narrow Feshbach resonances can be seen.

²Note that the position of the Feshbach resonance is not exactly where the molecular state of the bare closed channel crosses the incident scattering state. Instead, the coupling between the two channels creates new dressed states, and the actual position of the Feshbach resonance is where the *dressed* molecular state reaches the continuum.



Figure 2.2: The scattering length for Cs in the $|F = 3, m_F = 3\rangle$ state, given in units of the Bohr radius a_0 . The broad variation of the scattering length comes from a very broad *s*-wave Feshbach resonance at -11 G. On top of this several narrower *d*- and *g*-wave resonances at B = 11.0, 14.4, 15.0, 19.9, 48.0, 53.5, 112.8 and 131.1 G can be seen. The quantum numbers corresponding to the resonant molecular state are indicated with the notation (l, f, m_f) , where *l* is the orbital angular momentum, *f* is the internal angular momentum and m_f its projection on the magnetic field axis. Around 17.1 G, the scattering length has a zero crossing with a slope of $61 a_0/G$. Figure from [Chi04].

A very broad *s*-wave resonance centered at about -11 G leads to a slow variation of the scattering length, which makes it possible to tune the scattering length with a high precision. Especially interesting is the zero crossing at 17.1 G. Here, the scattering length varies with a slope of 61 G/ a_0 and, as will be demonstrated in section 4.2.3, it can be controlled with a precision better than 0.1 a_0 . A thourough discussion of Cs scattering properties and weakly bound molecular states can be found in [Chi01] and [Mar07b].

To be able to observe Feshbach resonances experimentally, the temperature has to be low enough such that the scattering cross section is not unitarity limited. This means that ka should not be much larger than one. For Cs atoms with a kinetic energy of $k_B \cdot 10 \,\mu\text{K}$, ka = 1 for $a \approx 270 \, a_0$.

Feshbach resonances offer the possibility to tune scattering length using an external magnetic field and are thus a very helpful tool for controlling atom-atom interactions. They can also be used to produce molecules out of an ultracold atomic gas by sweeping the magnetic field over the resonance, so called magneto-association [Her03, Reg03]. The unbound scattering state forms an avoided crossing with the bound molecular state in the closed channel and the magnetic field adiabatically converts the free atoms into molecules. This process is discussed in detail in several reviews [Köh06, Fer09].

2.1.2 Bose-Einstein condensation

An atomic gas of bosons behaves in different ways depending on its temperature. A qualitative picture of the different regimes is illustrated in figure 2.3. At high temperatures, the atoms in the gas behave as point-like particles. When the temperature is lowered, the atoms

2.1. Basic concepts



Figure 2.3: The behaviour of a gas of identical bosonic atoms at different temperatures. (A) At high temperatures, the gas can be treated as system of point-like particles. (B) For sufficiently low temperatures, the atoms must be described as wavepackets that scatter according to quantum mechanics. (C) A phase transition to a BEC occurs when the size of the atomic wavepackets is comparable to the mean distance between particles and the wavepackets start to overlap. (D) At zero temperature, all particles are in the same quantum state and can be described by a single macroscopic wave function. Adapted from [Ket99].

have to be described as quantum mechanical wave packets with an extent on the order of the de Broglie wavelength

$$\lambda_{\rm dB} = \frac{h}{\sqrt{2\pi m k_B T}},\tag{2.6}$$

where T is the temperature and m is the mass of the particle. The extent of the atomic wavepackets gets larger the further the temperature is lowered. At some point, the interatomic separation becomes comparable with size of the atomic wavepackets. The overlap of the atomic wavepackets can be quantified in terms of the phase-space density, defined as the density of the gas n multiplied by volume occupied by the wavepacket,

$$D = n\lambda_{\rm DB}^3. \tag{2.7}$$

When the phase-space density is on the order of unity, a phase transition will occur and the (bosonic) atoms form a Bose-Einstein condensate, where all atoms occupy the same quantum state. The atoms can then be described by a single macroscopic wave function. Note that to reach Bose-Einstein condensation, the gas must be sufficiently dilute that it does not become a liquid or a solid when being cooled. A BEC of atoms is in fact a metastable state, and will eventually decay through the formation of molecules.

There is a vast body of literature covering ultracold gases and Bose-Einstein condensation and I will here only review the parts of the subject that are relevant to the work presented in this thesis. Several textbooks [Pit03, Pet02] and review articles [Ket99, Dal99a, Cas01] provide further reading.

BEC of an ideal gas

Let us consider an ideal gas of bosons in thermal equilibrium with temperature T. The quantum state v will have a mean occupation number N_v given by the Bose distribution function

$$N_v = \frac{1}{e^{(\epsilon_v - \mu)/(k_B T)} - 1},$$
(2.8)

where ϵ_v is the energy of state v and μ is the chemical potential. For a fixed total number of particles N, the chemical potential is related to the temperature through the normalization

condition $N = \sum_{v} N_{v}$. The chemical potential is always smaller than the energy of the lowest state, $\mu < \epsilon_{0}$, since otherwise states with an energy lower than μ would have negative occupation numbers.

We can write the total number of particles as a sum of N_0 , the occupation number in the ground state, and a thermal component N_{th} , the number of particles in excited states.

$$N = N_0 + N_{\text{th}} = N_0 + \sum_{\nu=1}^{\infty} \frac{1}{e^{(\epsilon_\nu - \mu)/(k_B T)} - 1}.$$
(2.9)

For a fixed temperature, $N_{\text{th}}(\mu)$ varies smoothly and reaches a maximum for $\mu = \epsilon_0$. This means that the maximum number of particles in the thermal component is

$$N_{\text{th,max}} = \sum_{\nu=1}^{\infty} \frac{1}{e^{(\epsilon_{\nu} - \epsilon_{0})/(k_{B}T)} - 1}.$$
(2.10)

When the temperature is lowered, $N_{\text{th,max}}$ can be significantly lower than the total particle number N. This implies that a significant amount of particles must occupy the ground state, the signature of Bose-Einstein condensation. The temperature where $N_{\text{th,max}} = N$ and condensation starts is called the critical temperature T_c . If the thermal energy is much larger than the spacing between the energy levels, the sum in equation (2.10) can be replaced by an integral. For a harmonic trapping potential, the critical temperature can then be calculated to

$$T_{c} = \frac{\hbar\bar{\omega}}{k_{B}} \left(\frac{N}{\zeta(3)}\right)^{1/3} \approx 4.5 \frac{\bar{\omega}/2\pi}{100 \text{Hz}} N^{1/3} \,\text{nK},\tag{2.11}$$

where $\bar{\omega} = (\omega_x \omega_y \omega_z)^{1/3}$ is the geometrical average of trap frequencies and $\zeta(n)$ is the Riemann zeta function with $\zeta(3) \approx 1.2$. The critical temperature depends on both the atom number and the trap frequencies and is thus not constant during evaporative cooling. Note that for a large atom number, $(N/\zeta(3))^{1/3} \gg 1$ and the energy $k_B T_C$ corresponding to the critical temperature is much larger than the energy separation of the lowest trap levels. Still, the Bose statistics cause a macroscopic occupation of the ground state for temperatures below T_C .

It is often useful to monitor the progress towards Bose-Einstein condensation using the phase-space density of the gas instead of the temperature. The peak density of a classical gas in a harmonic trap can be derived from the Maxwell-Boltzmann distribution as

$$\hat{n} = N\bar{\omega}^3 \left(\frac{m}{2\pi k_B T}\right)^{3/2}.$$
(2.12)

Using equation (2.7), we can write the peak phase-space density in the trap on the form

$$D = N \left(\frac{\hbar\bar{\omega}}{k_B T}\right)^3.$$
(2.13)

From equation (2.11), we see that a temperature T_c would correspond to a critical phasespace density $D_c = \zeta(3) \approx 1.2$. However, equations (2.12) and (2.13) are derived for a classical gas and are not valid close to the critical temperature. It can be shown [Cas01] that in the limit $k_BT \ll \hbar\omega$, the critical phase-space density is

$$D_c = \zeta(3/2) \approx 2.6.$$
 (2.14)

Another useful quantity is the fraction of condensed atoms, which can be calculated from equations (2.10) and (2.11). The result is [Pit03]

$$\frac{N_0}{N} = 1 - \left(\frac{T}{T_c}\right)^3.$$
(2.15)

BEC of an interacting gas

A BEC of interacting particles in a trapping potential $V(\mathbf{r})$ can at zero temperature be treated in a mean-field approach [Gro61, Pit61, Pit03]. The BEC can then be described by a single macroscopic wave function $\Psi(\mathbf{r}, t)$ governed by the Gross-Pitaevskii equation

$$i\hbar\frac{\partial}{\partial t}\Psi(\boldsymbol{r},t) = \left[-\frac{\hbar^2}{2m}\nabla^2 + g|\Psi(\boldsymbol{r},t)|^2\right] + V(\boldsymbol{r})\right]\Psi(\boldsymbol{r},t),$$
(2.16)

where *g* is the interaction coupling constant given by equation (2.4) and $n(\mathbf{r}) = |\Psi(\mathbf{r})|^2$ is the density. This description is only valid in the dilute gas limit $n|a|^3 \ll 1$. Without interactions g = 0 and equation (2.16) reduces to the normal Schrödinger equation.

The stationary solution to equation (2.16) can be found by writing $\Psi(\mathbf{r}, t) = \Phi(\mathbf{r})e^{-i\mu t/\hbar}$, resulting in the time-independent Gross-Pitaevskii equation

$$\mu\Phi(\boldsymbol{r}) = \left[-\frac{\hbar^2}{2m}\nabla^2 + g|\Phi(\boldsymbol{r})|^2\right] + V(\boldsymbol{r})\right]\Phi(\boldsymbol{r}).$$
(2.17)

The energy of the system can be calculated from the wave function,

$$E = \int \left[\frac{\hbar^2}{2m} |\nabla\Psi|^2 + V(\mathbf{r})|\Psi|^2 + \frac{1}{2}g|\Psi|^4\right] d\mathbf{r} = E_{\rm kin} + E_{\rm pot} + E_{\rm int}.$$
 (2.18)

This expression contains three terms: E_{kin} is the quantum kinetic energy, often refered to as quantum pressure, E_{pot} is the potential energy of the system and E_{int} is the interaction energy or mean-field energy. From direct integration of the Gross-Pitaevskii equation (2.17), a relation between the chemical potential and the different energy terms can be derived:

$$\mu = \frac{E_{\rm kin} + E_{\rm pot} + 2E_{\rm int}}{N}.$$
(2.19)

An additional useful identity, known as the virial relation, is [Pit03]

$$2E_{\rm kin} - 2E_{\rm pot} + 3E_{\rm int} = 0.$$
 (2.20)

Thomas-Fermi approximation

In the limit of vanishing interactions, the solution to equation (2.17) with a harmonic potential is the harmonic oscillator ground state, the Gaussian wave function

$$\Phi(\mathbf{r}) = \sqrt{N} \left(\frac{1}{\sqrt{\pi}\sigma_{\text{ho}}}\right)^{3/2} e^{-\frac{r^2}{2\sigma_{\text{ho}}}},$$
(2.21)

19

where $\bar{\omega} = (\omega_x \omega_y \omega_z)^{1/3}$ is again the geometrical average of the trapping frequencies and $\sigma_{ho} = \sqrt{\hbar/(m\bar{\omega})}$ the associated harmonic oscillator length. However, in most cases the interactions cannot be neglected. A useful result comparing the interaction energy and the kinetic energy is

$$\frac{E_{\rm int}}{E_{\rm kin}} = \frac{N|a|}{\sigma_{\rm ho}}.$$
(2.22)

If the interaction energy dominates the kinetic energy, we can set the kinetic energy term in the Gross-Pitaevskii equation to zero. This is called the Thomas-Fermi approximation. The Gross-Pitaevskii equation will then reduce to

$$\mu = V(\mathbf{r}) + gn(\mathbf{r}). \tag{2.23}$$

The BEC density will arrange itself such that the interaction energy and the potential energy balance each other out and add up to a constant value, the chemical potential. The density distribution will therefore reflect the confining potential:

$$n(\mathbf{r}) = \begin{cases} \frac{\mu - V(\mathbf{r})}{g} & \text{where } \mu > V(\mathbf{r}), \\ 0 & \text{otherwise.} \end{cases}$$
(2.24)

For the common case of a harmonic confinement, the BEC density distribution has a parabolic shape, in contrast to the Gaussian density distribution of a thermal cloud or a non-interacting condensate.

It is straightforward to derive several useful physical quantities from the Thomas-Fermi solution. The peak density of the condensate is directly found in equation (2.24),

$$\hat{n} = \frac{\mu}{g} = \frac{\mu m}{4\pi\hbar^2 a}.$$
(2.25)

The total atom number is calculated by integrating the density, $N = \int n(\mathbf{r}) d\mathbf{r}$. This allows to deduce a relation between the chemical potential and the atom number in the case of a harmonic confinement,

$$\mu = \frac{1}{2}\hbar\bar{\omega} \left(\frac{15Na}{\sigma_{ho}}\right)^{2/5}.$$
(2.26)

The Thomas-Fermi-radii where the density becomes zero are calculated by setting $\mu = V(\mathbf{r})$,

$$R_{\text{TF},i} = \sqrt{\frac{2\mu}{m\omega_i^2}} = \sigma_{ho} \left(15N\frac{a}{\sigma_{ho}}\right)^{1/5} \frac{\bar{\omega}}{\omega_i} \quad i = x, y, z.$$
(2.27)

2.2 Experimental setup

2.2.1 The Cs atom

The experiments described in this thesis are carried out using Cs atoms. Cs is the heaviest stable member of the alkali metals and has only one stable isotope, ¹³³Cs. It is a solid at room temperature, with a melting point of 28°C. ¹³³Cs has a nuclear spin of 7/2, which together with the spin 1/2 from the single valence electron makes it a composite boson.



Figure 2.4: The energies of lowest ¹³³Cs electronic levels. The frequencies corresponding to the D1 and D2 transitions are $\omega_1 = 2\pi \cdot 335.1160487481(24)$ THz [Ger06] and $\omega_2 = 2\pi \cdot 351.72571850(11)$ THz [Ude00], respectively. The transition used for laser cooling has been marked.

The structure of the lowest electronic levels is shown in figure 2.4. The hyperfine splitting of the ground state ${}^{2}6S_{1/2}$ is the basis for the current definition of the second, with the splitting defined to be exactly 9.192631770 GHz. The $({}^{2}6S_{1/2} \rightarrow {}^{2}6P_{3/2})$ transition, referred to as the D₂ line, has a natural linewidth $\Gamma_{2} = 2\pi \cdot 5.23$ MHz. In this line, the transition $(F = 4, m_F = 4 \rightarrow F' = 5, m'_F = 5)$ is a closed transition that is suitable to use for laser cooling. Here, *F* and *F'* denote the hyperfine quantum number of the ground state and the excited state, respectively, and m_F and m'_F their projection on the magnetic field axis. The saturation intensity of this transition is $I_{sat} = 1.1 \text{ mW/cm}^2$. Another strong transition, not employed for laser cooling but important for optical dipole traps, is the $({}^{2}6S_{1/2} \rightarrow {}^{2}6P_{3/2})$ transition (D₁ line) which has a natural linewidth $\Gamma_{1} = 2\pi \cdot 4.58$ MHz.

¹³³Cs has a mass of 132.9 atomic mass units or $2.21 \cdot 10^{-25}$ kg. The recoil temperature, the temperature corresponding to an ensemble with a one-dimensional rms momentum of one photon recoil, is only 198 nK due to the large mass. This makes it possible to achieve very low temperatures by laser cooling only. The large mass also means that a Cs atom experiences a large potential gradient due to gravitation, $k_B \cdot 157 \,\mu\text{K/mm}$. A compilation of the physical and optical properties relevant to quantum optics experiments is [Ste02].

2.2.2 Vacuum system

Experiments with ultracold gases must be carried out in an ultra high vacuum (UHV), to minimize interactions with the room temperature lab environment. Our vacuum system, depicted in figure 2.2.2, is an evolution of the design used in the first generation Cs BEC setup in Innsbruck [Web03a, Her05]. It can be divided into five parts: An *oven* where Cs is heated to produce an atomic beam, the *oven pumping section* maintaining a pressure difference between the oven and the rest of the vacuum system, a long tube around which a *Zeeman slower* is





Figure 2.5: Overview of the vacuum system and the coils creating magnetic fields. The upper picture shows the vacuum chamber and the coils before the surrounding optics were installed. Also, the second part of the Zeeman slower had not yet been wound. Below is a CAD drawing of the setup with the different parts marked.

mounted, a *glass cell* where the experiments are carried out and the *main pumping section* providing the pumping capacity necessary to attain UHV. The whole system is mounted on an optical table with the centerline 300 mm above the table top.

Oven

The oven is the source for the Cs atoms that are cooled and trapped in the apparatus. It consists of a 64 mm diameter stainless steel tube with CF64 connectors. On one end there is a flange with 4 current feedthroughs and a viewport. 32 Cs dispensers (SAES Getters CS/NF/8/25 FT10+10) are mounted on Macor holders and connected to the feedthroughs in two independent circuits. When current flows through the dispensers, a pure Cs gas is released. We normally operate one of the circuits at 3.3 A, while the other is kept in reserve in case the first circuit should fail or run out of Cs.

A flange with a long nozzle, from which the atomic beam emerges, is attached to the other end of the oven. The nozzle, depicted in figure 2.6, is an aluminium piece consisting of two tubes with an opening milled in between. The first tube is 150 mm long with 3 mm inner diameter and extends far into the oven chamber. The choice of such a long tube allowed us to increase the diameter and thereby get a higher flux in the atomic beam while still maintaining good differential pumping. In addition, the geometry assures a well-collimated beam. The part of the nozzle residing outside the oven chamber has a 66 mm \times 14 mm square opening milled into it. The opening provides optical access to the atomic beam, for example for transversal cooling, and also allows for pumping between the tubes. After the opening, a 5 mm diameter, 96 mm long second tube serves as another differential pumping tube. Since the first and second tube are part of the same metal piece, their mutual alignment is automatically assured.

At room temperature, cesium is a solid with a melting point of 28°C and a vapor pressure of about 10^{-6} mbar [Ste02]. To produce an intense atomic beam, the pressure must be raised. The oven is thus heated by heating foil to 90°C, which corresponds to a Cs vapor pressure of $3 \cdot 10^{-4}$ mbar. Previous experience has shown that the viewport, which is used for diagnosis and alignment of the Zeeman slowing beam, can have its seal corroded by the cesium in the oven. To prevent accumulation of cesium on the seal, the viewport is kept 10°C warmer than the rest of the oven.



Figure 2.6: Oven nozzle. The long, narrow tube on the right extends into the oven chamber. The thicker part resides in the pumping section. The opening in the middle provides optical access to the atomic beam.

Oven pump section



Figure 2.7: Overview of the oven and the connected differential pumping section. Left: View from above. Right: View from the side.

The oven is connected to a section providing differential pumping, depicted in figure 2.7. The first part is a CF63 cube with viewports, providing optical access to the atomic beam. The cube has the same function as an ordinary 6-way cross but a more compact size. One side has a tee attached between cube and viewport with a 20 1/s ion getter pump (Varian Plus 20 StarCell) connected.

Having passed the cube, the atomic beam emerges through the differential pumping tube into a second pumping chamber, a 6-way cross. A titanium sublimation pump is connected to the lower flange of the cross. Titanium is a surface getter for almost all active gases and the surface area covered by the pump gives an estimated pumping speed of about 900 l/s for air. The pumping speed for Cs is not known, but the pumping efficiency for different gases is typically of the same order of magnitude as the pumping efficiency for air. A second 20 l/s ion getter pump pumps inert gases like noble gases or methane that are not absorbed by the titanium.

An ionization gauge (Varian UHV-24p) mounted on a tee measures the pressure in this region, about 10^{-10} mbar. On the opposite side of the tee is an angle valve (VAT series 54) that allows to connect an external turbomolecular pump and a roughing pump, to produce the initial vacuum needed for the titanium sublimation and ion getter pumps to work well.

A wobblestick connected to the upper flange of the 6-way cross serves as an atomic beam shutter. A small servo motor outside the vacuum can move the wobble stick in and out of the beam path. The oven section can be sealed off from the UHV part of the vacuum chamber with a pneumatic viton-sealed gate valve (VAT series 01). An electronic circuit monitors the pressure in the oven chamber and automatically closes the valve if the pressure rises above 10^{-8} mbar. This prevents a leak in the oven section from contaminating the UHV section.

Glass cell

The atoms are cooled and trapped in the main experiment chamber, a rectangular glass cell with inner dimensions 160 mm \times 65 mm \times 35 mm manufactured by Hellma GmbH. It is connected to the oven section through a 42 mm long, 16 mm diameter tube, where the



Figure 2.8: The main experiment chamber, a rectangular glass cell. The left, narrower glass tube connects to the Zeeman slower via a flexible bellows. The right, wider glass tube is connected to the UHV pumping section.

atomic beam is slowed down using the Zeeman slowing technique. The walls of the glass cell are made out of 6.5 mm thick glass plates polished to an optical surface quality of 15/15 scratch/dig. The glass material used, Vycor 7913, is transparent for light with wavelength between 300–2600 nm. The manufacturing process involves heating the glass plates to high temperatures, and an anti-reflection coating added to the glass plates before assembly would be destroyed. This means that the cell can only be coated after assembly, which limits any applied coating to the outside surfaces. Since it would thus not be possible to eliminate all reflections, we opted to not coat the outside either, making it possible to add laser beams at any wavelength in the future.

One end of the cell has a 28 mm inner diameter glass-to-metal-transition (Larson Electronic Glass) with a CF40 flange which connects to the Zeeman slower tube via a flexible edge-welded bellow (ComVAT). The bellow relieves mechanical stress which would otherwise break the fragile glass chamber. The other end of the cell has a larger 58 mm inner diameter glass-to-metal-transition with a CF63 flange, which is fastened to a 5-way cross providing pumping. The large diameter was chosen to assure a high conductance to the 5-way cross, achieving high pumping speeds. Both flanges are made out of low magnetic permeability SAE type 316 stainless steel.

The mounting of the glass cell to the rest of the vacuum system is a delicate procedure. The two flanges of the cell cannot both be rigidly mounted, since there would be a high risk of imposing mechanical stress, especially considering that the vacuum system experiences large temperature changes during bakeout. On the other hand, letting one flange float freely is not an option since its weight would impose a large torque on the cell. We therefore decided to let the flange rest on a support residing on a spring. The spring was adjusted such that it exactly cancels the gravitational force acting on the flange.

To seal the CF flanges we used annealed copper gaskets, which are softer than normal copper gaskets. This minimizes the stress applied to the cell while mounting. The bolts connecting the flanges were tightened sequentially one quarter turn at a time. To assure that the flanges were well aligned and not askew, their gap between them was periodically measured with shims.

UHV pumping section

A 5-way cross with attached pumps maintains ultra-high vacuum in the glass cell. A titanium sublimation pump coats a surface corresponding to a pumping speed of 4000 l/s for air. The pumping speed in the glass cell is however limited by the aperture connecting the pump section to the glass cell, with a conductance of about 100 l/s for the Cs gas coming from the oven. Inert gases are pumped by a 55 l/s ion getter pump (Varian Plus 55 Star-

Cell). A viewport provides optical access for the Zeeman slower beam, an all-metal angle valve (VAT series 57) allows connection to external pumps, and an ionization gauge (Varian UHV-24p) measures the pressure in the chamber.

The pressure measured by the ionization gauge is off-scale, indicating a pressure of less than 10^{-11} mbar. It is difficult to tell what the pressure in the glass cell is, but the lifetime of an atomic sample in a magnetic quadrupole trap is above 60 s, which is more than enough for our purposes.

Cleaning and baking

In a UHV system, the base pressure is determined by the flow of gases into the chamber and the available pumping capability. If the system is leak tight, only a slight amount of He diffusing through glass surfaces will come into the system from the outside. Instead, most of the gas load will come from internal sources. Apart from the atoms in the atomic beam, these sources are mainly outgassing from surface contamination and diffusion of impurities from the stainless steel chamber walls. The speed with which the impurities diffuse out of the chamber walls increases exponentially with temperature, and to attain a good base pressure one has to heat the chamber while the whole system is under vacuum.

Before assembly, all stainless steel parts were cleaned with acetone in an ultrasonic bath to remove surface contamination. The system was then assembled, evacuated and baked in two stages. First, the system was assembled with viewports and the glass cell replaced with stainless steel components. The setup was heated to 350° C for one week while under vacuum, pumped by a turbomolecular pump. The pressure was then about 10^{-7} mbar. After cooling down to room temperature and flashing the titanium sublimation pumps, a base pressure below 10^{-11} mbar was measured. The system was then flooded with argon gas, viewports and glass cell were connected and the Zeeman slower coil mounted around the Zeeman slower tube. With all glass parts mounted, a second bakeout was performed. To make sure that the glass cell would not break due to thermal expansion, the system was slowly heated by 5°C/h until 200°C was reached. Great care was taken to isolate the system to prevent large thermal gradients. The system was baked for 10 days and then slowly cooled down to room temperature again.

In the early stages of the experiment, a leak in the oven viewport led us to break vacuum in the oven section. After the leak was repaired and the system closed off again, the chamber has now been under vacuum for several years without any maintenance except for periodic activation of the titanium sublimation pumps.

2.2.3 Magnetic fields

It is crucial to have both fast and precise control over the magnetic field in the experiment region. We have installed a number of coils serving different purposes. A pair of Helmholtz coils provides a large vertically oriented homogenous field to tune the Cs scattering length. Two coils in anti-Helmholtz configuration create a quadrupole field, providing a field gradient which is used to counteract gravity. A large cage creates a smaller homogenous field in arbitrary direction to compensate for the earth magnetic field and provide the field necessary for Raman sideband cooling. Finally, a set of solenoids create the field for the Zeeman slower. The coil design is in many ways inspired by the designs used in the Innsbruck GOST experiment [Eng06].

Helmholtz and quadrupole coils

Working with ¹³³Cs has the advantage that only moderate magnetic field strengths are needed, both to tune the scattering length and to access narrow Feshbach resonances where molecules can be created. Additionally, our optical trapping approach means that no magnetic trap is need, avoiding the requirement of high magnetic field gradients. The goal when designing the main coils that produce a homogenous field and a quadrupole field was therefore not to optimize for maximum field strength, but rather to maximize the optical access, keep the coil inductance small and the water cooling simple.

The homogenous field is created by a pair of coils with corotating currents. Their radius R and mutual distance D was chosen such that R = D, the Helmholtz configuration, which maximizes the field homogeneity. The coils were designed to provide a magnetic field of up to 150 G, allowing us to tune the scattering length up to 1600 a_0 .

Another coil pair creates a quadrupole field. Here, the geometry has several optima – for a given mutual distance D, the field gradient is maximal if R = D and the homogeneity is maximal for $R \approx 0.58D$. The requirement to be able to create a gradient strong enough to compensate gravity for both atoms and Feshbach molecules, about 60 G/cm, could be easily met and we choose a radius $R \approx 0.8D$, a compromise between a homogenous gradient and large gradient strength allowing good optical access.

The coils are made out of a single layer of flat 1 mm \times 8 mm copper wire. They are mounted on water-cooled aluminium plates, which dissipates the heat from the coils. Every winding has good thermal contact with the plate, which assures adequate cooling. The cooling plates are milled out of a 6 mm thick aluminium sheet. They are slit in order to avoid eddy currents which reduce the speed with which the magnetic field can be switched. To further reduce the amount of metal between the coils and the trapped atoms, additional grooves have been milled into the plate. The disc is cooled by water flowing through a copper tube along the periphery.

The individual coils were wound by hand around a plastic template, which was affixed into a lathe. To ensure that the coils held together, a thin layer of epoxy was applied to the first and last five windings. We used the epoxy Emerson & Cumming EccoBond 285 / Catalyst 9, with a good thermal conductivity and low thermal expansion coefficient which had worked well in a neighboring lab [Eng06]. This epoxy was also used to bond the coils to the cooling plate. It was important to use very thin layers of epoxy to achieve good heat transfer, applying just enough to fill any pockets of air between the coil and the cooling plate. When the cooling plate is well cooled, a continuous current of 70 A through the quadrupole coils, giving 60 G/cm field gradient, leads to a temperature rise of about 25°C,.

The coils are mounted around the glass cell as depicted in figure 2.9. There is one cooling plate above the glass cell and one below, separated by spacers made of Tufnol, a machinable non-conducting material. Each cooling plate has an inner and an outer coil attached to each side, for a total of eight coils. The four outer ones, with 11 windings each, are connected in series to effectively form a pair of coils in Helmholtz configuration and provide a homogenous field. The four inner coils, each with 15 windings, are connected to form a pair of coils with counterrotating current and create the quadrupole field. The most important properties of the coils are summarized in table 2.1. The magnetic field was calibrated using microwave spectroscopy as described in section 2.2.8, or by measuring the position of the known Feshbach resonances [Chi04]. The measured fields were found to agree very well with the expected values.



Figure 2.9: Left: Drawing depicting how the Helmholtz and quadrupole coils are mounted around the glass cell, viewed from the side. **Right:** Picture of the Helmholtz and quadrupole coils, viewed from above.

The magnetic field response to a sudden change in coil current was probed using the imaging beam. Starting with a large field, the current through the coils was switched off within microseconds and the field at the position of the atoms was probed by measuring the resonance frequency of the imaging beam. From this we could estimate the 1/e decay time to be about 200 µs, considerably faster than in the first generation Cs BEC setup [Her05] which used a stainless steel vacuum chamber.

	Helmholtz coils	Quadrupole coils
Radius	77 mm	60.5 mm
Coil separation	76 mm	76 mm
Number of windings	2×22	2×30
Resistance	$74~\mathrm{m}\Omega$	$65 \mathrm{m}\Omega$
Inductance	290 µH	250 µH
Field strength (calculated)	2.59 G/A	0.85 (G/cm)/A
Field strength (measured)	2.58 G/A	0.83 (G/cm)/A

Table 2.1: Properties of the main Helmholtz and quadrupole coils.

Current control

The Helmholtz and quadrupole coils are powered by two programmable switched-mode power supplies (Delta Elektronika SM60-100 and SM30-100). They provide a maximum current of 100 A and a maximum voltage of 60 V and 30 V, respectively. Fast ramping speeds are achieved using a feedback loop, where the current is measured with a sensitive current transducer (DanFysik UltraStab 867-200I) and fed back to a PID circuit which controls the voltage of the power supplies.

Upward ramps of the current *I* are limited by the current rise speed dI/dt = V/L, where *V* is the applied voltage and *L* the coil inductance. If the maximum voltage is applied, the current rises with 200 A/ms for the Helmholtz coils and 120 A/ms for the gradient coils. This assures that the current can be ramped up to any desired value within less than 500 µs. Downward ramps are limited by the fact that the power supplies are unipolar. Since the

voltage required to drive the coils in steady-state is only a few V, the best one can do is set the voltage to zero, which causes the current to decay exponentially with the time constant L/R, where R is the resistance of the circuit. We have therefore installed 0.5 Ω external resistors in series with the coils, which for both coil pairs improves the time constant from about 4 ms to 500 µs.

Even with external resistors installed, the absolute ramp speed is limited at low currents because the current decays exponentially when the voltage is set to zero. This problem was overcome by installing fast high-power diodes in the circuit, which only conduct current above a threshold voltage. The zero of the voltage scale is therefore effectively shifted, and applying zero voltage at the power supplies is equivalent to applying a negative voltage across the coils. This makes fast ramping possible even at low currents.

The current can be quickly switched off using MOSFET transistors. The energy stored in the coils is then dissipated in varistors connected in parallel with the transistors. It takes about 10 μ s to switch the current off.



Compensation coils

Figure 2.10: Geometry of the compensation coils, which can create a homogenous field of about 1 G in an arbitrary direction.

A large cage consisting of three square coil pairs creates a homogenous magnetic field of about 1 G in arbitrary direction. Its geometry is illustrated in figure 2.10. The primary use of the compensation coils is to cancel static fields like the earth magnetic field and stray fields from lab equipment like ion getter pumps, but it is also used to provide the small field needed for Raman sideband cooling and to compensate for the field of the Zeeman slower. The large size of the cage was chosen for several reasons; the coils are far away from the main coils, which minimizes the inductive coupling. This avoids induced currents in the compensation circuits when the main coils are quickly ramped. Large coils also produce more homogenous fields. Finally, it is convenient to have the coil frame far away from the experiment chamber, thereby not removing any optical access.

The measured resistance, inductance and field strength of the compensation coils is shown in table 2.2. Self-made linear regulated power supplies can send up to 2A current through

Direction	Resistance	Inductance	Measured field strength
x	3.5Ω	12 mH	225 mG/A
у	$1.5 \ \Omega$	1.1 mH	309 mG/A
Z	$1.6 \ \Omega$	2.1 mH	426 mG/A

the coils. A new current can typically be set in about 1 ms.

2.2.4 Zeeman slower

The atomic beam coming out of the oven is decelerated in a Zeeman slower [Phi82, Met99]. The atoms are slowed by scattering photons from a counter-propagating laser beam. A carefully designed magnetic field introduces a Zeeman shift that compensates for the change in Doppler shift due to the decreasing velocity. To achieve a constant deceleration, the magnetic field profile should have a square-root shape.

To be able to capture atoms with high velocities, the Zeeman slowing region has to be long. On the other hand, the atomic beam diverges during the propagation through the Zeeman slower, and making the Zeeman slower longer does not automatically ensure a larger flux of slow atoms through the MOT capture volume. For light atoms, the divergence is mainly due to a spread of the transverse velocity distribution due to the scattering of many photons. Cs is heavy and has a low recoil momentum, hence the transverse heating is not as large and the main factor determining the divergence is mainly due to the initial transverse velocity component of the atoms when emerging from the oven, which becomes important at the end of the slower where the longitudinal velocity is small. The large beam divergence at the end of the slower also implies that it is important to keep the distance between the MOT and the Zeeman slowing section short. We therefore opted for a decreasing-field Zeeman slower where the field of the Zeeman slowing coils seamlessly merges with the MOT quadrupole field. In this way, the atoms can be decelerated all the way to the MOT.

Simulations of the slowing process indicated an optimal Zeeman slower length of about 70 cm. The Zeeman slower solenoid has to be split in two parts to accommodate a post supporting the flange of the UHV glass cell. The first part resides around the long tube connecting the oven section to the glass cell and consists of a 80 mm diameter, 480 mm long brass tube around which a 1x2.5 mm insulated flatband copper wire was wound. After winding one layer, the wire was fixated with epoxy (UHU Plus Endfest 300) before moving on to wind the next layer. The first four layers were wound along the whole length of the tube, with 178 windings per layer. Thereafter an additional 12 layers were wound, with the number of windings in each layer was gradually reduced to produce the necessary magnetic field profile. A measurement of the magnetic field with a Hall probe showed excellent agreement with the calculated profile.

The second of the Zeeman slower coils resides around the glass-metal-transition of the glass cell. This part had to be wound in place after the vacuum setup was assembled. Two halves of a plastic tube were mounted in a rotatable way around the glass-to-metal-transition, such that the coil could be wound by simply rotating the tube. The coil has a total of 7 layers of windings in a conical shape, with 12 extra windings close to the edge of the coil to help

Table 2.2: Measured properties of the compensation coils. *x*-, *y*- and *z*-directions are defined in figure 5.10.

bridge the gap to the first coil.

Figure 2.11 shows the calculated magnetic field profile. The second of the Zeeman slower coils creates a field of about 2 G at the position of the MOT. This field is compensated by the compensation coils. The coils are operated at a current of 2 A except for the four inner layers of the first coil, which are run at 4.4 A. The two Zeeman slower coils dissipate a total power of 40 W, which means that no water cooling is necessary.



Figure 2.11: Left: Schematic drawing of the Zeeman slower section with the two conical Zeeman coils. **Right:** The axial magnetic field of the Zeeman slower and the quadrupole coil. The red curve is the total field and the blue curves show the contributions from the individual coil layers.

2.2.5 Magnetic levitation

Cs is the heaviest of the stable alkali atoms and the gravitational force causes Cs atoms to experience a large potential gradient of $k_B \cdot 157 \,\mu\text{K/mm}$. Since the dipole traps we use have trap depths on the order of tens of μ K, it is immediately clear that larger dipole traps will not be able to support the atoms against gravity. We therefore employ the magnetic levitation technique [Web03a, Her05] and apply a vertical magnetic field gradient, causing an upward force that counteracts the gravitational force.

For atoms in $|F = 3, m_F = 3\rangle$, the potential due to a magnetic field is

$$U_{\rm mag} = -\mu(B)B \approx \mu^{(1)}B - \mu^{(2)}B^2, \qquad (2.28)$$

with

$$\mu^{(1)} = h \cdot 1.50283 \text{ MHz/G} = 0.7522 \,\mu_B,$$

$$\mu^{(2)} = h \cdot 93.5 \text{ MHz/G}^2. \tag{2.29}$$

where $\mu_B = 9.27400915 \cdot 10^{-28}$ J/G is the Bohr magneton. The second term in equation (2.28) is due to the quadratic Zeeman effect. The gravitational force can be counteracted by applying a magnetic field gradient such that $\partial U_{mag}/\partial z = mg$. The gradient needed to support the atoms against gravity is

$$\frac{\partial B}{\partial z} = 31.02 \text{ G/cm} - 0.0055 \text{ G/cm} \cdot B_0/\text{G}, \qquad (2.30)$$

and is slightly dependent on the field amplitude B_0 due to the quadratic Zeeman effect.



Figure 2.12: Equipotential lines of the magnetic field resulting from a quadrupole field and a homogenous offset field oriented in the *z*-direction. From [Her05].

Unfortunately, applying a vertical gradient automatically leads to outward horizontal forces. This can be understood by drawing the equipotential lines of a quadrupole field, as seen in figure 2.12. The atoms, residing above the magnetic field minimum, experience a force perpendicular to the equipotential lines. It is clear that the closer the atoms are to the quadrupole center, the stronger the curvature of the potential becomes. Neglecting the quadratic Zeeman effect, the magnetic potential in the radial direction ρ is [Eng06]

$$U_{\rm mag}(\rho) = \mu^{(1)} B - m\omega_{\rm lev}^2 \rho^2, \tag{2.31}$$

a repulsive harmonic potential with (anti-)trapping frequency

$$\omega_{\rm lev} = g \sqrt{\frac{m}{4\mu^{(1)}B_0}}.$$
(2.32)

For a field $B_0 = 17$ G, the antitrapping frequency is $2\pi \cdot 3.4$ Hz.

2.2.6 Diode laser system

For laser cooling and absorption imaging, narrow-band laser light with a wavelength close to the Cs D_2 line at 852 nm is needed. This is provided by a set of home-built diode lasers. These have been described in great detail in two diploma theses [Unt05, Fli06] and I will here only summarize the main points of the setup.

The laser system has to provide the following frequencies (the hyperfine quantum numbers of the ground state and the excited state of the D_2 line are denoted F and F', respectively):

- − MOT cooling light, tunable between 70-0 MHz red-detuned from the main laser cooling transition ($F = 4 \rightarrow F' = 5$).
- MOT repumper light to pump off-resonantly scattered atoms from F = 3 back into F = 4, resonant with $(F = 3 \rightarrow F' = 3)$.
- Zeeman slower light, 50 MHz red-detuned from ($F = 4 \rightarrow F' = 5$).
- Zeeman slower repumper, 50 MHz red-detuned from ($F = 3 \rightarrow F' = 3$).
- The optical lattice for Raman sideband cooling, resonant with $(F = 4 \rightarrow F' = 4)$.

- Polarizer beam, 9 MHz blue-detuned from $(F = 3 \rightarrow F' = 2)$.
- Imaging beam, resonant with $(F = 4 \rightarrow F' = 5)$.

A grating-stabilized diode laser in Littrow configuration serves as master laser. It is locked 320 MHz red-detuned to the ($F = 4 \rightarrow F' = 5$) transition using an error signal obtained by modulation transfer spectroscopy. With this locking scheme a linewidth of about 10 kHz is achieved. In addition to its use as a frequency reference, this laser also provides a few mW of light for absorption imaging.

The light for the Zeeman slower comes from a slave laser, a laser diode which is locked using the injection lock technique. A few mW of seed light is split off from the master laser, frequency shifted using acousto-optical modulators (AOMs) and injected into the laser cavity. If the laser current and temperature is right, the slave laser will inherit the spectral properties of the seed light without the need for any locking electronics. For diagnosis, saturated absorption spectroscopy is implemented.

Two further slave lasers, also injection-locked, provide light both for the MOT and the optical lattice used for Raman sideband cooling. Two different seed beams with the frequencies needed for the MOT and the Raman lattice are overlapped on a 50/50 beam splitter, coupled into an optical fibre and injected into the slave laser cavity. The lasers can then be switched to the frequency needed within microseconds by simply turning on the corresponding seed beam. The light output by the lasers can be directed to either the MOT or the Raman lattice using AOMs.

The laser light addressing transitions from the F = 3 hyperfine state is about 9 GHz detuned from the master laser, which is too far detuned to easily reference to the master laser. Instead, we use a second grating-stabilized diode laser. It is built similar to the master laser and locked 125 MHz red-detuned from the ($F = 3 \rightarrow F' = 2$) frequency, reaching a linewidth of 30 kHz. From this laser the light for Zeeman slower repumping, MOT repumping and and the polarizer beam used in Raman sideband cooling is derived using AOMs as frequency shifters.

The diode laser setup is built on an optical table separate from the experiment table. The light is transported to the experiment using single-mode optical fibers. This assures excellent beam quality and decouples the beam pointing at the experiment from the details of the laser setup, which significantly facilitates maintenance. Except for the Zeeman slower light, every laser beam has at least one AOM in its beam path, which makes it possible to switch the light or control the intensity with microsecond precision. In addition, shutters are installed before each fiber to mechanically block any remaining stray light. The setup has proved to be robust, and since the installation of an air conditioning unit with enough power to keep a stable temperature in the lab, the experiment can run for many hours with all lasers remaining locked.

2.2.7 Experiment control

BEC creation takes several seconds. On the other hand, many of the processes in the experiment require laser intensities, laser frequencies, magnetic fields, RF and microwave signals to be controlled with microsecond precision. An experiment control system thus needs to provide precise and reproducible timing on both short and long time scales. In addition, many digital and analog outputs channels are needed.



Figure 2.13: Schematic of the hardware for experiment control.

We have chosen a combination of hardware and software developed by F. Schreck and T. Meyrath³, illustrated in figure 2.13. The central hardware is a computer equipped with a buffered fast digital I/O card (National Instruments PCI-6533), which provides 32 digital outputs that can be switched at up to 2 MHz. 25 of these outputs are used to create a simple data bus which can communicate with different output devices. On the bus, 8 bits are used to select which device to address, 16 bits are used for sending data to the device and one bit is used as a clock. The bus is distributed to different parts of the lab by means of flat ribbon cables. Anywhere on the flat ribbon cable, an output device can be connected. We currently have three different types of output devices in use:

- **Digital output board** Provides 16 digital TTL outputs, which can all be switched simultaneously.
- Analog output board Provides 8 analog outputs with 16 bit resolution and an output range of ± 10 V. Since the bus is 16 data bits wide, only one output can be changed every clock cycle. The settling time for an output port is 10 µs.
- Digital frequency generator This device uses an Analog Devices AD9852 Direct Digital Synthesizer (DDS) to create radio frequency signals. The DDS chip is capable of creating frequencies between 0-135 MHz with 48 bit resolution, corresponding to 0.5 µHz. The amplitude can be controlled with 12 bit precision.

Designs exist for other devices like analog input boards and mechanical shutter drivers, but are not employed in the current setup.

There are several advantages to this hardware concept. It is easy and cheap (around 50 EUR for 16 digital outputs, 150 EUR for 8 analog outputs and 150 EUR for a frequency generator) to add new output boards anywhere in the lab. This means that one can put output ports near the devices that are to be controlled, reducing the need to clutter the lab with long cables, as would be the case when using one large central output board. The ability to use shorter cables is not only convenient, it also reduces the chance of the analog lines picking up stray signals. We currently have 80 digital outputs, 24 analog outputs and 15 digital frequency generators available, distributed over different parts of the lab.

The system clock is provided by a signal generator (SRS DS345) phase locked to a stable 10 MHz reference. Many environmental parameters, like for example magnetic fields, fluctu-

³The system is documented at http://george.ph.utexas.edu/control/index.html

ate with the 50 Hz frequency of the AC power line, and it is therefore desirable to keep every experimental run synchronized with the power line. This is accomplished by a home-built electronics device which, on receiving a TTL signal, stops the clock signal from reaching the main digital I/O card until the power line has reached a predetermined phase. When no clock signal is present, the I/O card suspends the output of data from its buffer, which is resumed when the clock signal comes back. This way of synchronizing the experiment with the power line has the advantage that it can be done anytime during the timing sequence. A line sync can thus be done at the beginning of the experimental run, and if a magnetic field sensitive operation needs to be performed several seconds later in the timing sequence when the AC line phase may have drifted compared to the much more stable system clock, another line sync can be triggered.

The hardware is controlled by a software frontend written in C++. The software is capable of programming arbitrary waveforms and allows flexible flow control. A set of experimental parameters can be defined, which can be varied through a simple user interface. After every experimental run, the parameters used are logged to the hard drive. Measurement series can be created where one or more parameters are varied in a nested loop. It is also possible to put several measurement series in a queue so that a long series of different measurements can be executed one after another, allowing the apparatus to run unattended.

2.2.8 Detection and diagnosis

Imaging

The density distribution of the atomic sample is determined using absorption imaging [Ket99]. A laser beam resonant with the ($F = 4 \rightarrow F' = 5$) transition illuminates the atoms and the resulting shadow is imaged on a CCD chip. The amount of light absorbed is related to the column density of the atom cloud, and by comparing the shadow image to a reference image taken with no atoms present, the density distribution and atom number can be determined. Due to the photons absorbed by the sample, absorption imaging is destructive and new experimental run is needed for every picture.

We use a camera (Apogee Alta U32+, CCD chip Kodak KAF-3200ME) with a pixel size of 6.8 μ m × 6.8 μ m, 14 bit resolution and an active area of 2184 × 1472 pixels. The image of the atom cloud is collected by a 25 mm diameter, *f* = 80 mm achromatic doublet optimized for diffraction limited performance (Melles Griot 01LAI009) at a distance of 80 mm from the trap center and focused onto the CCD chip by a 25 mm diameter, *f* = 175 mm achromat (Melles Griot 01LAO167). This lens arrangement has a theoretical resolution limit of 3.1 μ m and provides a magnification of 2.2. The lenses are mounted in a long lens tube that is directly connected the camera. This provides automatic alignment and shields the CCD chip from stray light. The magnification of the lens system was determined by measuring the position of a freely falling atom cloud and fitting the result with a parabola, and we found that one pixel corresponds to 3.1 μ m × 3.1 μ m, in agreement with the calculated value.

The absorption images are always taken with the magnetic fields turned off. The MOT repumping light is turned on 100 μ s before the absorption picture is taken, to pump the atoms into the F = 4 hyperfine state. The imaging beam is then pulsed on for 150 μ s. The images are read out from the camera to a computer and visualized and processed with a Matlab program. Gaussian or bimodal fits can be applied to the density profiles, from which useful quantities like atom number, Gaussian width of a thermal cloud or Thomas-Fermi radius of

a BEC are inferred. The program can plot these quantities as a function of different experimental parameters and automatically apply different fit functions to the data. The real-time visualization is very helpful, as it allows immediate assessment of ongoing measurements.

A triggerable video camera (Pulnix TM6-EX) mounted at approximately 30° to the vertical provides information about the atomic sample from a second perspective. This camera is mainly used to take fluorescence pictures, where the MOT light is switched on for a short time and the fluorescence of the atoms is imaged. Although the limited resolution and signalto-noise ratio of the camera is not useful for quantitative analysis, it has proven very useful when aligning the different dipole traps.

Thermometry

The temperature of the atomic cloud is measured using time-of-flight measurements. The trapping potential is shut off and an absorption picture is taken a time *t* later. The experiment is repeated for different expansion times *t* and the Gaussian width $\sigma(t)$ of the freely expanding cloud is measured. The temperature *T* can then be determined from the relation

$$\sigma(t) = \sqrt{\sigma_0^2 + \frac{k_B T}{m} t^2},\tag{2.33}$$

where σ_0 is the size of the cloud at time t = 0 and m is the atom mass.

For temperatures in the nK range, it is necessary to use long expansion times get a reliable temperature measurement. The magnetic levitation field is therefore kept on during the expansion. The transversal force due to the levitation field causes the cloud to spread in the horizontal direction and equation (2.33) is only valid for the vertical expansion. For the horizontal expansion, one has to use the relation [Her05]

$$\sigma(t) = \sqrt{\sigma_0^2 \cosh^2(\omega_{\text{lev}} t) + \frac{k_B T}{m} \frac{\sinh^2(\omega_{\text{lev}} t)}{\omega_{\text{lev}}^2}},$$
(2.34)

where ω_{lev} is the levitation anti-trapping frequency defined in equation (2.32).

Microwave spectroscopy

The application of microwave fields at around 9.2 GHz can drive transitions between the Cs lower and upper hyperfine ground states F = 3 and F = 4. Every state experiences a different Zeeman shift, therefore the frequency needed to drive a transition between two states is dependent on the value of magnetic field. Microwave spectroscopy can thus be used as a sensitive tool to measure the magnetic field.

A frequency generator capable of producing frequencies up to 3 GHz (Rohde & Schwarz SMIQ03B) was readily available in the lab, which prompted us to adapt a very simple setup to create the 9.2 GHz radiation needed to drive transitions between the hyperfine states. The frequency generator, referenced to an accurate 10 MHz signal derived from a commercial frequency standard (SRS PRS10), is set to output a 2.3 GHz signal. The signal is sent to a frequency doubler (Eclipse D1550L), an amplifier (Mini-Circuits ZX60-6013E) and a second frequency doubler (Eclipse D2060L). This creates a 9.2 GHz signal, which is fed to a power amplifier (ALC APA0612-35-30) and then sent to an outcoupler situated above the experiment chamber at a distance of about 10 cm from the trap center.


Figure 2.14: Measurement of the magnetic field using microwave spectroscopy. The atom loss from the trap was measured as a function of the detuning from the hyperfine splitting, 9.9192631770 GHz. A fit to the data yields a center frequency of 42.063(4) MHz, which corresponds to a magnetic field of 17.143(2) G.

An example of a magnetic field calibration is shown in figure 2.14. A trapped cloud of atoms in the $|F = 3, m_F = 3\rangle$ state was irradiated with a microwave field tuned to populate the $|F = 4, m_F = 4\rangle$ state. Inelastic collisions lead to heating and result in a loss of atoms from the trap. In this case, the measured atom loss after 200 ms is maximal for a detuning of 42.063(4) MHz from the 9.9192631770 GHz hyperfine splitting.

Including the quadratic Zeeman shift, the shift of the energy for the magnetic sublevel m_F in the hyperfine states F = 3 and F = 4 can be calculated to [Eng06]

$$E_{F=3,m_F}/\text{Hz} = -h \cdot 350943.7875m_F \cdot (B/\text{G}) - 13.355(16 - m_F^2) \cdot (B/\text{G})^2,$$
(2.35)

$$E_{F=4,m_F}/\text{Hz} = +h \cdot 349820.3875m_F \cdot (B/\text{G}) + 13.355(16 - m_F^2) \cdot (B/\text{G})^2.$$
(2.36)

These expressions are valid for low magnetic fields. The error is about 10^{-6} for 10 G and about 0.1 % for 100 G. For the transition ($|F = 3, m_F = 3\rangle \rightarrow |F = 4, m_F = 4\rangle$), the magnetic field corresponding to a detuning ν from the hyperfine splitting can be calculated from equations (2.35) and (2.36):

$$B(\nu) = 0.407811 \frac{\text{G}}{\text{MHz}} \nu \left(1 - \frac{\nu}{64319 \text{ MHz}} \right).$$
(2.37)

For the measurement in figure 2.14, the magnetic field is then determined to be 17.143(2) G.

2.3 The path to BEC

Cs has a low recoil energy due to its large mass, which makes it possible to achieve low temperatures with laser cooling. Together with the availability of laser diodes from the optical telecommunication industry at around 850 nm, the wavelength of the Cs D₂ line, this made Cs one of the main candidates in the quest to achieve a quantum degenerate gas [Tie92, Mon93]. However, attempts to evaporatively cool Cs in a magnetic trap were marred by spin-changing collisions, both for Cs in the $|F = 4, m_F = 4\rangle$ state [Söd98, Arl98] and in the $|F = 3, m_F = -3\rangle$ state [GO98, Hop00, Tho03].

A way to get around this problem is to trap the lowest Zeeman sublevel, $|F = 3, m_F = 3\rangle$, where spin-changing collisions are energetically forbidden at low temperatures. This state

2. A new Cs BEC apparatus

is a high-field seeker and cannot be trapped in a magnetic trap. Instead an optical trap has to be used, an approach followed in experiments in Stanford [Vul99], Berkeley [Han01] and two different experiments in Innsbruck, one surface trapping experiment [Ham02] and one using a conventional dipole trap [Web03c]. With the problem of two-body inelastic collisions solved, a new enemy turned up in the form an unusually large three-body recombination rate, leading to trap loss and heating [Web03c]. The Innsbruck experiment still managed to reach BEC using forced evaporation in the dipole trap in a narrow magnetic field window around 21 G [Web03b]. It was later found that three-body recombination is reduced at this magnetic field, an effect associated with few-body Efimov physics [Kra06].

Later, Cs BEC has been achieved in the Innsbruck surface trap [Ryc04], in the apparatus described in this thesis and, recently, at the University of Chicago [Hun08].

We follow the path to BEC pioneered in the Innsbruck experiment, summarized in table 2.3. Atoms are laser cooled in two steps, first in a *magneto-optical trap* and then via 3D *Raman sideband cooling* to below 1 μ K. A large, shallow *crossed-beam dipole trap* is then loaded, with a magnetic field gradient balancing gravity for the trapped $|F = 3, m_F = 3\rangle$ state. This magnetic levitation also automatically assures perfect spin polarization. The large-volume dipole trap can catch many of the laser cooled atoms but is not suitable for evaporation due to its low trap frequencies. Instead, a second, tightly focused dipole trap, the *dimple trap*, is loaded. *Forced evaporation* is then done by lowering the laser power of the dimple trap while keeping the magnetic field at the three-body recombination minimum at 21 G. Using these techniques, we can produce a BEC with up to $2 \cdot 10^5$ atoms every 10 s.

	Density (cm^{-3})	Temperature	PSD	Atom number
Oven	$\approx 10^7$	360 K	$pprox 10^{-21}$	N/A
MOT	$pprox 10^{10}$	$pprox 70 \ \mu K$	$\approx 5 \cdot 10^{-8}$	$3 \cdot 10^8$
Compressed MOT	$\approx 4 \cdot 10^{10}$	$pprox 40 \ \mu K$	$\approx 5 \cdot 10^{-7}$	$2 \cdot 10^8$
Raman sideband cooling	$2\cdot 10^{10}$	650 nK	$1 \cdot 10^{-4}$	$6\cdot 10^7$
Large dipole trap	$1\cdot 10^{11}$	1.1 µK	$4\cdot 10^{-4}$	$1\cdot 10^7$
Dimple trap	$4\cdot 10^{13}$	1.1 μK	$1\cdot 10^{-1}$	$1.5\cdot 10^6$

Table 2.3: Summary of temperatures and densities for the different cooling stages. PSD denotes phase-space density. The exact values can vary due to changes in the lab environment, shown are typical good values. A well aligned setup produces a BEC with about $2 \cdot 10^5$ atoms.

2.3.1 Magneto-optical trap

The magneto-optical trap (MOT) has become the standard technique to capture atoms and cool them down to μ K temperatures [Met99]. Our MOT is created by three pairs of counter-propagating laser beams with 10 mm waist. Two beam pairs are located in the horizontal plane, at an angle of 75° to each other. For these beams 65 mW of light from a diode laser is available. A second diode laser provides 70 mW power for the third beam pair, which is oriented at a slight angle from the vertical. The intensities are about 10 mW/cm² for the horizontal beams and 22 mW/cm² for the vertical beams. Additionally, 10 mW of repumping light is overlapped with the vertical beams. The magnetic confinement is created by the main quadrupole coils.

The MOT captures atoms out of the Zeeman slowed atomic beam. 75 mW of light reddetuned 50 MHz from the laser cooling transition ($F = 4 \rightarrow F' = 5$) is available for the slowing together with 8 mW of repumping power. The slowing laser beam is expanded to a beam waist of about 2 cm and enters the vacuum system through a viewport at the end of the UHV pumping section (see section 2.2.2). The beam is focused down such that its focus is slightly behind the nozzle of the oven. The beam can exit through a second viewport at the end of the oven, which is very helpful when aligning the beam. If the beam can pass cleanly through the long oven nozzle, it is well aligned with the atomic beam.

We found that MOT loading works best with a magnetic field gradient of 4.5 G/cm and the MOT light 10 MHz red-detuned from the cooling transition. The initial loading rate from the Zeeman slower is about $2 \cdot 10^8$ atoms/s and after 2 s of loading, approximately $3 \cdot 10^8$ atoms are collected. The whole atom cloud does not fit within the limited field-of-view of our imaging setup when expanded such that it is not optically dense, and these numbers should be considered rough estimates. The flux of slow atoms could be improved by implementing transverse cooling of the atomic beam emerging from the oven. Optical access to implement such cooling is available, and the rather slow velocity of the Cs atoms out of the oven would allow for a long interaction time with the cooling beams. A rough estimate indicates that it should be possible to increase the loading rate by about an order of magnitude, but since the current loading rate is satisfactory for our current purposes, we have not implemented this yet.

After the MOT is loaded, it is compressed [DeP00] by increasing the magnetic field gradient to 25 G/cm over 25 ms, and then holding that gradient for 25 ms. At the same time, the MOT detuning is ramped to 65 MHz. In this way the density is increased, which considerably increases the efficiency when the laser cooled cloud is later transfered to the optical dipole trap. After compression, we typically end up with $2 \cdot 10^8$ atoms at a temperature of about 40 μ K. We have found that the temperature of the compressed MOT is not very important for the efficiency of the next laser cooling step, Raman sideband cooling, and therefore made no effort to reach lower temperatures. To prepare for Raman sideband cooling, the repumper light is shut off 2 ms before the compressed MOT is turned off, which makes the atoms fall into the $|F = 3\rangle$ state.

2.3.2 3D Raman sideband cooling

In most BEC experiments, atoms are loaded out of a MOT into a magnetic trap or optical dipole trap, where evaporative cooling is performed. We have chosen to add a further laser cooling step, 3D Raman sideband cooling, before loading the dipole trap. This allows us to reach temperatures well below 1 μ K and simultaneously polarize the atoms into the state $|F = 3, m_F = 3\rangle$, which is the state we want to trap later. I will here describe the main points of this cooling stage, a more detailed characterization can be found elsewhere [Fli06].

Our setup follows the approach outlined in [Ker00, Tre01] and is illustrated in figure 2.15. The atoms from the MOT are trapped in a three-dimensional optical lattice. Due to their relatively high temperature, they will initially populate highly excited vibrational trap levels. A magnetic field is applied such that the Zeeman splitting of states with $\Delta m_F = 1$ is equal to the vibrational level spacing in the lattice. This means that the state $|\nu, m_F\rangle$ with vibrational quantum number ν and hyperfine sublevel m_F is degenerate with the state $|\nu - 1, m_F - 1\rangle$. In addition to trap the atoms, the lattice light also drives Raman transitions between these states. Simultaneously, the polarizer, a mainly σ^+ -polarized beam, provides optical pumping

2. A new Cs BEC apparatus

to $P_{3/2}$, F' = 2. In the Lamb-Dicke regime, ν is conserved during the pumping and the whole process therefore leads to rapid loss of vibrational quanta. The state $|\nu = 0, m_F = 2\rangle$, which is not resonant with σ^+ -polarized light, is depopulated by a weak π -polarized component of the polarizer beam. In the end, all atoms are transferred to $|\nu = 0, m_F = 3\rangle$, a dark state. The atoms are then released into free space by adiabatically ramping down the lattice, which reduces the free-space temperature [Kas95].



Figure 2.15: Left: Schematic of the cooling process for Raman sideband cooling. Atoms in vibrationally excited states are cooled by degenerate Raman transitions (double-sided arrows) and optical pumping with a strong σ^+ component and a weak π component. Figure from [DD04].

Right: Five laser beams are needed for 3D Raman sideband cooling, one circularly polarized beam for optical pumping (blue) and four linearly polarized beams for the lattice (red).

There is 120 mW of laser power available for the lattice, which is formed by four beams with 1.5 mm $1/e^2$ -radius. Geometry, relative intensities and polarizations were originally chosen as in [Tre01] and have later been empirically optimized to maximize atom number and minimize temperature. The lattice light comes from two injection-locked diode lasers and is resonant with the ($F = 4 \rightarrow F' = 4$) transition, thus being 9.2 GHz detuned with respect to the resonance for atoms in $|F = 3\rangle$. The lattice therefore acts as a repumper for atoms off-resonantly scattered into $|F = 4\rangle$. The relatively small lattice detuning means that cooling is essential to trap atoms – the heating rate was measured to about 0.2 μ K/ms with the polarizer beam was turned off, leading to a rapid loss of atoms from the lattice.

The polarizer is 9 MHz blue-detuned from the ($F = 3 \rightarrow F' = 2$) transition. The detuning is necessary to compensate for the light shift from the lattice potential. The beam has 120 µW power, a 1.5 mm $1/e^2$ -radius and the light is circularly polarized.

The magnetic field is created by the compensation coils and is oriented along the vertical, parallel to the polarizer beam. If the polarizer had perfect circular polarization, only σ^+ -transitions would be driven. Any ellipticity gives rise to σ^- -transitions, which leads to heating. The magnetic field is therefore slightly tilted by adding a small component in the horizontal direction. The elliptical polarization is then converted into σ^+ light with a small amount of π light, which is exactly what is needed. We found that a linear ramp of the magnetic field amplitude during the cooling process brought a significant gain in the number of trapped and cooled atoms. This is probably due to the finite size of the lattice beams which leads to different trap frequencies at the edge of the lattice compared to the center, requiring different magnetic fields for efficient cooling.

In only 3 ms of cooling, up to $6 \cdot 10^7$ atoms are trapped, polarized into $|F = 3, m_F = 3\rangle$ and cooled to about 650 nK. This corresponds to a phase-space density of about $1 \cdot 10^{-4}$. The phase-space density has increased by three orders of magnitude in a few ms, which shows the remarkable efficiency of this cooling scheme. These are good starting conditions for loading the optical dipole trap.

2.3.3 Large volume dipole trap - reservoir

Because of the low temperature of the atom cloud after Raman sideband cooling, only a shallow trap is needed to trap the atomic sample. This makes it possible to use large beams with attainable laser power. A large-volume dipole trap has the advantage that it can trap a large part of the laser-cooled atoms. The gravitational force is compensated with the magnetic levitation technique described in section 2.2.5. Gravitation is only balanced for atoms in the $|F = 3, m_F = 3\rangle$ state and atoms in other states are lost from trap, assuring that the sample is always perfectly spin-polarized.

Optical dipole traps

A light field induces an electrical dipole moment in an atom, which in turn interacts with the electric field of the light. The result is a potential proportional to the intensity of the light, which can be used to trap atoms. For an atomic transition with resonance frequency ω_0 and natural linewidth Γ , the potential depth $V(\mathbf{r})$ and the photon scattering rate Γ_{sc} are [Gri00]

$$V_{\rm dip}(\boldsymbol{r}) = -\frac{3\pi c^2 \Gamma}{2\omega_0^3} \left(\frac{1}{\omega_0 - \omega} + \frac{1}{\omega_0 + \omega}\right) I(\boldsymbol{r})$$
(2.38)

$$\Gamma_{\rm sc}(\boldsymbol{r}) = \frac{3\pi c^2 \Gamma \omega^3}{2\hbar \omega_0^6} \left(\frac{1}{\omega_0 - \omega} + \frac{1}{\omega_0 + \omega}\right)^2 I(\boldsymbol{r})$$
(2.39)

This expression is valid when the detuning $|\omega - \omega_0|$ is much larger than the linewidth Γ and the photon scattering is low ($\Gamma_{sc} \ll \Gamma$).

In our case, the two atomic resonances corresponding to the Cs D_1 and D_2 lines have to be taken into account. ω_0 then has to be replaced by an effective transition frequency [Gri00]

$$\omega_{\rm eff} = \frac{1}{3}\omega_1 + \frac{2}{3}\omega_2 = 2\pi \cdot 2.56 \cdot 10^{14} \,\,{\rm Hz}.$$
(2.40)

Similarly, Γ has to be replaced by an effective linewidth

$$\Gamma_{\rm eff} = \frac{1}{3}\Gamma_1 + \frac{2}{3}\Gamma_2 = 2\pi \cdot 5.00 \text{ MHz.}$$
(2.41)

We use dipole traps with 1064 nm light. For this wavelength, the potential depth and the scattering rate is

$$V_{\rm dip} = k_B \cdot 2.6 \cdot 10^{-6} \,\mu \text{K} \cdot I / (\text{mW/cm}^2), \qquad (2.42)$$

$$\Gamma_{\rm sc} = 1.6 \cdot 10^{-8} \,{\rm s}^{-1} \cdot I/({\rm mW/cm^2}) = 0.0061 \,{\rm s}^{-1}/(V_{\rm dip}/{\rm \mu K}). \tag{2.43}$$

41

2. A new Cs BEC apparatus

The intensity profile of a Gaussian laser beam with power P is

$$I(\rho, z) = \frac{2P}{\pi w^2(z)} e^{-2\rho^2/w^2(z)} \quad \text{with } w(z) = w_0 \sqrt{1 + \left(\frac{z}{z_R}\right)^2}, \tag{2.44}$$

where z is the coordinate along the beam axis, ρ the coordinate in radial direction and w(z) the beam waist. $z_R = \pi w_0^2 / \lambda$ is the Rayleigh range and w_0 the beam waist at the focus of the beam. Even for beam waists w_0 as small as 100 µm, the Rayleigh range is 3 cm. The confinement in the axial direction is therefore very weak and can safely be neglected for our purposes.

The trap depth can be calculated from equations (2.38) and (2.44). Our large-volume dipole trap has beams with a waist of 500 μ m, which leads to a trap depth

$$V_0 = k_B \cdot 0.65 \,\mu\text{K/W}.$$
 (2.45)

To achieve confinement in all directions, two crossed beams can be overlapped. In this case, the total trap depth is given by the beam with the lowest trap depth.

Technical setup

The light for the large-volume dipole trap is generated by a 100 W Ytterbium fiber laser (IPG Photonics YLR-LP-100) with 1071 nm wavelength. The laser is running on many longitudinal modes, causing the emitted light to have a 3 nm spectral width. The beam coming out of the fiber is specified to be spatially single-mode with $M^2 < 1.1$, but we have found that the beam profile becomes severely distorted at high output power. We thus normally run the laser at about 40 W, where the beam shape is still fine.

The light has a stable linear polarization, and the laser beam is split in two beam paths with equal power using a half-wave plate and a polarizing beam splitter. When we first set up the trap, we used AOMs to switch the light on and off. Several lenses had to be introduced into the beam path to focus the light through the AOM and the small beam waist needed led to very high intensities. Although great care was taken to keep the lab environment clean, we found that dust burned on the mirrors and the AOM crystals, leading to drifts in the beam pointing and distortion of the beam profile. We have therefore opted for an alternative, simple solution that minimizes the number of optical elements in the beam path. The light is switched on an off using mechanical shutters, where a mirror mounted on a motor can be moved into the beam path to deflect the laser beam into a water-cooled beam dump. The switching can be done with ms precision, considerably less precise than an AOM but enough for our purposes. Care was taken to keep the beam waist large, especially where the beam hits an optical surface. To minimize reflections, the two beams enter the experiment chamber at Brewster's angle, which for the Vycor glass used in our glass cell is 56° to the normal of the surface. Both beams propagate in the horizontal plane and cross at an angle of 68°. Telescopes where one lens is on a translation stage allow for easy and reproducible adjustment of the beam waists in the experiment region in the range $300-1200 \,\mu\text{m}$. We found that the transfer to the dipole trap worked best with a beam waist of 500 μ m.

Trap loading

To transfer the laser-cooled atoms into the dipole trap, the dipole trap is turned on during the last second of laser cooling. When the cloud is released from the Raman lattice, the magnetic

levitation gradient is ramped up. The gradient is first ramped to a value of 60 G/cm, more than needed to counteract gravity, within 700 μ s. It is then relaxed to 31.3 G/cm within 350 μ s, exactly compensating for gravity. This deliberate "overshoot" compensates for the finite ramping speed and cancels the downward velocity acquired while the levitation was ramped up. Simultaneously to the levitation gradient, the magnetic offset field is ramped up to 129 G, corresponding to a scattering length of 1540 a_0 .

The transfer into the dipole trap heats the sample to about 3 μ K. The temperature is then reduced by plain evaporation, causing the atom number to be reduced to $1 \cdot 10^7$ and the temperature to 1.1 μ K. The initial heating can be attributed to imperfect mode-matching between the dipole trap and the Raman-cooled cloud. In the dipole trap, a 650 nK sample in thermal equilibrium would have an 1/e-radius of about 200 μ m. The Raman-cooled cloud has a 1/e-radius of about 800 μ m and is badly mode-matched to the dipole trap. It therefore has an excess of potential energy that is converted into kinetic energy and heats the sample.

2.3.4 Tightly focused dipole trap - dimple

A large-volume dipole trap is required to ensure a high transfer efficiency from the laser cooling stage. It is however not a good trap for forced evaporation because of the low trap frequencies and the low density of the atomic sample. Additionally, the broad-band trapping light we use cannot be used to trap Cs_2 molecules because the trap light would induce transitions to excited states. To overcome these problems we use a second "dimple" trap, a more tightly focused crossed dipole trap with narrow-band light.

When the strength of a trapping potential is adiabatically changed, an increase in density of the trapped sample is accompanied by an increase in temperature, and the phase-space density remains constant. A non-adiabatic change will even decrease the phase-space density. However, as first demonstrated in [Pin97], changing the shape of the trapping potential can locally increase the phase-space density. If we have a large trap filled with cold atoms and introduce a narrow dimple in the potential, elastic collisions will load atoms into the dimple. As long as the number of atoms in the dimple is small compared to the reservoir of atoms in the large trap, the temperature will remain almost constant. The density in the dimple will be increased by the Boltzmann factor $\exp(U/k_BT)$, where U is the depth of the dimple potential. The local phase-space density will thus be increased, at the prize of a lower atom number.

Technical setup

A fiber amplifier, described in section 3.2.1, produces light with 1064 nm wavelength and 1 kHz linewidth for the dimple trap. From this light, two beams are derived. Each beam is sent through an AOM and is then coupled into a polarization-maintaining single-mode optical fiber. The light intensity after the fiber is monitored using a photodiode measuring the light leaked through a dielectric mirror. The electronic circuit reading out the photodiodes features a logarithmic amplifier, which allows to precisely measure the intensity over several orders of magnitude. The photodiode signal is fed to PID circuit, which stabilizes the light intensity using the AOM. This ensures a stable trap depth over several orders of magnitude unaffected by fluctuations in laser intensity and fiber transmittance.

One of the dimple beams propagates in the horizontal plane, at an angle of 45° to the glass cell. It is focused down to a beam waist of 40 µm by a lens with focal length 300 mm.

2. A new Cs BEC apparatus

The lens is mounted on a translation stage to allow precise adjustment of the focus. The other dimple beam enters the chamber vertically. It has a beam waist of 150 μ m in the experiment region. We chose to let one of the dimple beams propagate vertically because it can be used to augment the very weak radial confinement of the 1D optical lattice (described in section 3.2.1).

Trap loading

The dimple is loaded by raising the power to 80 mW in the tightly focused horizontal beam and 550 mW in the vertical beam in a 1 s linear ramp. The horizontal beam creates a 7.5 μ K deep potential with $2\pi \cdot 170$ Hz trap frequency and the potential created by the vertical beam is 4 μ K deep with $2\pi \cdot 100$ Hz trap frequency. It is desirable to have a large scattering length while the dimple is loaded, assuring enough collisions to keep the low density sample in the reservoir in thermal equilibrium. On the other hand, a large scattering length enhances losses due to three-body recombination in the dimple, where the density is high. We found a scattering length of 450 a_0 to be a good compromise. After the dimple is ramped up, we lower the scattering length to 300 a_0 .

After 300 ms, one reservoir beam is shut off and the atoms residing in the reservoir escape. The second reservoir beam is initially kept on, since the vertically propagating dimple beam is not strong enough to provide confinement on its own. About $1.5 \cdot 10^6$ atoms remain in the dimple trap at a temperature of about 1.1 μ K. The density is about $4 \cdot 10^{13}$ cm⁻¹ and the phase-space density is 10^{-1} . These are excellent conditions for evaporative cooling.

2.3.5 Evaporation and Bose-Einstein condensation

We start forced evaporative cooling by ramping down the power of the dimple beams. The remaining reservoir beam is turned off when the power in the horizontal dimple beam is 40 mW. The evaporation is most efficient when the scattering length is set to 210 a_0 , where the elastic collision rate is large enough to allow for fast evaporation and the three-body recombination rate has a local minimum [Kra06]⁴. We use a piecewise exponential ramp over 6 s, where the power is in the horizontal dimple beam is reduced from 80 mW to 1 mW. The power in the vertical dimple beam is typically reduced from 550 mW to 10 mW.

At the end of the ramp, a nearly pure BEC with up to $2 \cdot 10^5$ atoms remains in the trap. The trap depth at the end of the ramp is 280 nK and the trap frequencies $2\pi \cdot (20, 20, 28)$ Hz. For a scattering length of 210 a_0 , this corresponds to Thomas-Fermi radii of (14,14,10) µm and a peak density of $5 \cdot 10^{13}$ cm⁻¹.

An example of the tunability of the BEC interaction energy is shown in figure 2.17. The BEC is released from the trap and imaged after different times of free expansion. At the time of release from the trap, the scattering length is set to a new value, 10 a_0 and 100 a_0 , respectively. The interaction energy of the BEC, which is proportional to the scattering length (see section 2.1.2) is converted to kinetic energy and determines the rate of expansion. The difference in expansion rate can be clearly seen. Note also that the expansion rate is much larger in the horizontal direction than in the vertical direction. This is due to the horizontal anti-trapping caused by the magnetic levitation, as explained in section 2.2.5. When the

⁴When starting with a low atom number in the dimple, it was necessary to set the scattering length to a higher value towards the end of the evaporation ramp to keep the elastic collision rate large enough.



Figure 2.16: Density distribution of the atomic sample with the evaporation ramp ending at different laser power of the dimple trap, imaged after 50 ms levitated time-of-flight. The phase transition from a thermal cloud with a Gaussian density distribution (**left**) to a BEC with a much narrower Thomas-Fermi distribution (**right**) is clearly seen.

scattering length is set close to zero, the expansion energy becomes very small [Web03b]. We have measured expansion energies as low as $k_B \cdot 200$ pK.

2. A new Cs BEC apparatus



Figure 2.17: Levitated expansion of a BEC with the scattering length set to $10 a_0$ (left) and $100 a_0$ (right), respectively.

3

A BEC in an optical lattice

I will in this chapter review the basic theoretical tools needed to describe a BEC in a 1D optical lattice and describe the experimental setup and techniques used in our apparatus. For further reading there are several reviews of both experiments and theory of ultracold atoms in optical lattices, for example [Mor06] and [Blo08].

3.1 Theory

3.1.1 Optical lattice potential

Two superimposed counter-propagating narrow-band laser beams will interfere and create a standing wave pattern. If the laser beams are far detuned from the atomic resonance and have radially symmetric Gaussian beam profiles, the resulting periodic potential at the beam focus takes the form

$$V(r,z) = -V_0 e^{-2r^2/w_0^2} \cos^2(kz),$$
(3.1)

where $k = 2\pi/\lambda$ is the wave vector of the laser light, w_0 is the $1/e^2$ -radius (waist) of the beam, V_0 is the lattice depth, z is the direction of beam propagation and r is the radial distance from the beam center. It has been has assumed that the confinement in z-direction is negligible, e.g. $z \ll z_R$, where z_R is the Rayleigh length. An easy way to create two counter-propagating laser beams is to simply retro-reflect a single beam, as illustrated in figure 3.1.

Due to constructive interference, the lattice depth V_0 is four times larger than the depth of a non-retroreflected dipole trap with the same intensity. It can be calculated using equation (2.38),

$$V_0 = -\frac{8P}{\pi w_0^2} \frac{3\pi c^2 \Gamma}{2\omega_0^3} \left(\frac{1}{\omega_0 - \omega} + \frac{1}{\omega_0 + \omega} \right)$$
(3.2)

It is often useful to specify the lattice depth as $V_0 = sE_R$, where $E_R = \hbar^2 k^2/2m$ is the recoil energy imparted by one lattice photon on an atom with mass m. For ¹³³Cs atoms and 1064 nm lattice light, $1 E_R = k_B \cdot 64$ nK = $h \cdot 1325$ Hz.

In the center of the trap, the potential can be well approximated by a sum of a harmonic radial confinement and a homogenous lattice potential:

$$V(r,z) = \frac{1}{2}m\omega_r^2 r^2 + V_0 \cos^2(kz).$$
(3.3)

47



Figure 3.1: The simplest way to create a standing wave is to retro-reflect a Gaussian laser beam. Illustration from [Win07a].

Here the radial trapping frequency is given by

$$\omega_r = \sqrt{\frac{4V_0}{mw_0^2}}.\tag{3.4}$$

In some cases, it is useful to approximate the individual lattice wells with harmonic traps. For a given lattice depth V_0 , the trap frequency along the lattice direction z is

$$\omega_{\text{lat}} = \sqrt{V_0 \frac{2k^2}{m}} = 2\sqrt{s}\omega_{\text{rec}},\tag{3.5}$$

where $\omega_{\rm rec} = \hbar k^2 / 2m$ is the recoil frequency.

3.1.2 Bloch states and band structure

The motion of a particle in a potential V(x) is governed by the Hamiltonian

$$\hat{H} = \frac{\hat{p}^2}{2m} + V(x).$$
(3.6)

In free space, where V(x) = 0, the set of plane waves $|p\rangle$ carrying momentum p form a complete basis of solutions to the Schrödinger equation,

$$\hat{H}|p\rangle = E(p)|p\rangle$$
 where $E(p) = \frac{p^2}{2m}$. (3.7)

When V(x) is a periodic potential, the equivalent role is played by the *Bloch states* $|\phi_q\rangle$. Similar to as plane waves, these states are delocalized in space. To calculate these states, we have to seek the eigenstates to the Schrödinger equation

$$\hat{H} |\phi_q\rangle = E_n(q) |\phi_q\rangle \tag{3.8}$$

where the potential V(x) in our case comes from an optical lattice and is sinusoidal,

$$V(x) = V_0 \cos^2(kx) = V_0 \left(\frac{1}{2} + \frac{1}{4}e^{i2kx} + \frac{1}{4}e^{-i2kx}\right).$$
(3.9)

48

3.1. Theory

Since multiplying a wave function with a factor e^{ikx} is equal to adding a momentum $\hbar k$, the potential can also be written in operator form as

$$\hat{V} = \frac{1}{4} V_0 \left(2 + \hat{T}_{+2\hbar k} + \hat{T}_{-2\hbar k} \right), \qquad (3.10)$$

where $\hat{T}_{\pm 2\hbar k}$ denotes an operator adding the momentum $\pm 2\hbar k$ to the wave function. This can be understood in a simple picture; an atom can absorb a photon from one lattice beam, which is followed by stimulated emission into the counter-propagating beam, causing a total transfer of two recoil momenta $\hbar k$ from the lattice to the atom.

It is therefore natural to seek the eigenstates $|\phi_q\rangle$ of the Hamiltonian as a superposition of plane waves with momentum spacing of $2\hbar k$. This leads to the following ansatz:

$$|\phi_q\rangle = \sum_l c_l(q) |q+l \cdot 2\hbar k\rangle.$$
(3.11)

The "central momentum" q is called *quasimomentum* and can without loss of generality be restricted to the first Brillouin zone $[-\hbar k, \hbar k]$. The coefficients $c_l(q)$ can be calculated by inserting the ansatz into the Schrödinger equation (3.8). The kinetic energy term in the Hamiltonian is then

$$\frac{\hat{p}^2}{2m} |\phi_q\rangle = \sum_l c_l(q) \frac{(q+l\cdot 2\hbar k)^2}{2m} |q+l\cdot 2\hbar k\rangle$$
(3.12)

$$=\sum_{l}c_{l}(q)\left(\frac{q}{\hbar k}+2l\right)^{2}E_{R}\left|q+l\cdot2\hbar k\right\rangle.$$
(3.13)

Using equation (3.10), the potential energy term is

$$\frac{1}{4}V_0\left(\hat{T}_{+2\hbar k} + \hat{T}_{-2\hbar k} + 2\right)|\phi_q\rangle =
= \frac{1}{4}V_0\sum_l c_l(q)\left[|q + (l+1)\cdot 2\hbar k\rangle + |q + (l-1)\cdot 2\hbar k\rangle + 2|q\cdot 2\hbar k\rangle\right]
= \frac{1}{4}V_0\sum_l \left[c_{l-1}(q) + 2c_l(q) + c_{l+1}(q)\right]|q + l\cdot 2\hbar k\rangle.$$
(3.14)

The Schrödinger equation can now be written in matrix form as

$$\sum_{l'} H_{l,l'} c_{l'}^n(q) = E_n(q) c_l^n(q) \quad \text{with} \quad H_{l,l'} = \begin{cases} (q/\hbar k + 2l)^2 E_R + V_0/2 & \text{for } l = l' \\ V_0/4 & \text{for } l - l' = \pm 1 \\ 0 & \text{elsewhere.} \end{cases}$$
(3.15)

Equation (3.15) is an eigenvalue problem which can be solved numerically by truncating the Hamiltonian matrix $H_{l,l'}$; for lattice depths $V_0 \leq 20 E_R$ it is sufficient to take $|l| \leq 5$ into account if only the lowest energy bands are considered. One can thus calculate the eigenenergies $E_n(q)$ for Bloch states with different quasimomenta q. The wave function corresponding to a certain Bloch state can be determined from the coefficients $c_l^n(q)$ using equation (3.11).

The dispersion relation $E_n(q)$ is shown in figure 3.2 for different lattice depths. In the limit of vanishing lattice depth, the Bloch states are equal to plane waves and the dispersion



Figure 3.2: Illustration of Bloch bands. The energies of the Bloch states plotted against their quasimomentum for different lattice depths. With increasing lattice depth, the lowest band flattens and the band gap approaches the level spacing $\hbar \omega_{lat}$ of a harmonic oscillator.

relation is equal to the dispersion relation in free space limited to the first Brillouin zone. With a lattice present, the range of possible energies is split up into energy bands called *Bloch bands*. As the lattice gets deeper, the gap between the bands increases and the width of the bands decreases. For very deep lattices, the gap between the lowest bands approaches the level spacing of the harmonic oscillator $\hbar \omega_{lat}$.

In figure 3.3, Bloch functions are shown in both position space and momentum space. In position space, one can see that the Bloch functions are delocalized over the entire lattice. As could be expected, the probability density has maxima centered over the lattice sites. In momentum space, the concept of Bloch states as a superposition of discrete momentum states is easily seen.

3.1.3 Wannier states

The Bloch states form an orthogonal set of eigenstates to the Schrödinger equation delocalized over the entire lattice. For a shallow lattice, this a good choice of basis to work with. When working with deep lattices, it is often more convenient to work with a set of states localized to a single lattice site. There exists such an orthogonal set of states, called *Wannier states*. A Wannier state is constructed as a specific superposition of all Bloch states in one band [Ash76],

$$|w^{(n)}\rangle = \mathcal{N}^{-1/2} \sum_{q} |\phi_q^n\rangle, \qquad (3.16)$$

where $|w^{(n)}\rangle$ denotes the Wannier state corresponding to a particle in the n^{th} band and \mathcal{N} is a normalization constant¹. The sum is carried out over all q belonging to the first Brillouin

¹Note that equation (3.16) does not define a unique set of Wannier states, since the Bloch states $|\phi_q\rangle$ are arbitrary up to a complex phase. For the lowest band, choosing the Bloch states such that the coefficients $c_l(q)$ are real and positive gives a Wannier function that is real, symmetric and falls off exponentially. There exists only one such Wannier function per band, called a maximally localized Wannier function [Koh59]. Throughout this work, this particular set of Wannier functions will be used.

3.1. Theory



Figure 3.3: Density distribution of the Bloch states for quasimomenta 0, 0.5 and 1 $\hbar k$ in position space (**upper**) and momentum space (**lower**). In the lower plot, it can be clearly seen how the Bloch states are a superposition of momentum states spaced 2 $\hbar k$ apart. The dotted line marks the Wannier function (see next section).



Figure 3.4: Left: Density distribution of Wannier functions for different lattice depths *s*. The inset shows the sidelobes close to next neighbouring lattice site **Center, right:** Comparison of the Wannier function and the harmonic oscillator ground state (3.17) for 5 E_R and 15 E_R deep lattices. Note the discrepancy in amplitude at the neighbouring lattice site.

zone. Using equation (3.16), one gets a Wannier state localized over lattice site 0. The Wannier function at lattice site j can be inferred from the function at lattice site 0 by a simple translation $w_j^{(n)}(z) = w_0^{(n)}(z - jd)$, where $d = \pi/k$ is the lattice site spacing. In momentum space this translates to $w_j^{(n)}(p) = e^{-ipjd/\hbar}w_0^{(n)}(p)$.

Note that the Wannier states are *not* eigenstates to the lattice potential. Their virtue is that they are localized states that form a basis for a single Bloch band. Since most of the following discussion will concern the lowest Bloch band, we will drop the band index n and refer to the Wannier functions in the lowest band as $w_j(z)$.

For sufficiently deep lattices, the lattice wells can be seen as isolated harmonic traps and the Wannier functions approach the harmonic oscillator ground state

$$w_{ho}(x) = \frac{1}{\pi^{1/4} \sigma_{\text{lat}}^{1/2}} e^{-\frac{z^2}{2\sigma_{\text{lat}}^2}},$$
(3.17)

where σ_{lat} is the harmonic oscillator length of the lattice well,

$$\sigma_{\text{lat}} = \frac{d}{\pi s^{1/4}}.$$
(3.18)

The Gaussian approximation can be made more accurate by, instead of using σ_{ho} as the Gaussian width, seeking the width that minimizes the energy of the system. One finds [Cri02] that a Gaussian with a width σ_c satisfying the condition

$$e^{-(\sigma_c/\sigma_{\rm lat})^2/\sqrt{s}} = (\sigma_c/\sigma_{\rm lat})^{-4}$$
 (3.19)

is the best approximation.

As can be seen in figure 3.4, there is one important difference between Wannier functions and the harmonic oscillator ground state; Wannier functions have sidelobes at neighbouring lattice sites. Therefore, even for deep lattices, the wave function overlap with neighbouring lattice sites is severely underestimated in a Gaussian approximation.

Conversion between quasimomentum and momentum space

Let us assume that we have a wave function characterized by a quasimomentum distribution f(q). How does this wave function look in momentum space? We can write the wave function as a sum of Bloch functions

$$\Psi_f(p) = \int f(q)\phi_q(p)dq.$$
(3.20)

A Bloch function can be written as a sum of Wannier functions multiplied by a phase factor [Har04],

$$\phi_q(p) = \frac{1}{\sqrt{\hbar k}} \sum_{l=-\infty}^{\infty} w_n(p) e^{ildq/\hbar}$$
(3.21)

$$=\frac{1}{\sqrt{\hbar k}}w_0(p)\sum_{l=-\infty}^{\infty}e^{ild(q-p)/\hbar},$$
(3.22)

where we have used the relation $w_j(p) = w_0(p)e^{-ipjd/\hbar}$. The sum over l can be transformed using the identity $\sum_{l=-\infty}^{\infty} e^{ils/S} = 2\pi S \sum_{l=-\infty}^{\infty} \delta(s+l \cdot 2\pi S)$, and we can then write

$$\phi_q(p) = \sqrt{\hbar k} w_0(p) \sum_{l=-\infty}^{\infty} \delta(q - p + l \cdot 2\hbar k).$$
(3.23)

We can now rewrite equation (3.20) and deduce that a quasimomentum wave function f(q) in momentum space is equal to

$$\Psi_f(p) = \int \left(f(q)\sqrt{\hbar k} \, w_0(p) \sum_{l=-\infty}^{\infty} \delta(p-q+l\cdot 2\hbar k) \right) dq \tag{3.24}$$

$$=\sqrt{\hbar k}\,w_0(p)\tilde{f}(q).\tag{3.25}$$

The wave function in momentum space is thus simply $\tilde{f}(q) = \sum_{l=-\infty}^{\infty} (p+l \cdot 2\hbar k)$, the quasimomentum wave function periodically repeated, multiplied by a Wannier function envelope and a normalization constant.

3.1.4 Effective 1D equation

Let us now consider a system consisting of a three-dimensional BEC trapped in a harmonic trap with a superimposed one-dimensional optical lattice. In the weakly interacting regime, the dynamics of the system is governed by the Gross-Pitaevskii equation

$$i\hbar\frac{\partial}{\partial t}\Psi(\boldsymbol{r},t) = \left[-\frac{\hbar^2}{2m}\nabla^2 + g|\Psi(\boldsymbol{r},t)|^2 + V_{\text{ext}}(\boldsymbol{r})\right]\Psi(\boldsymbol{r},t)$$
(3.26)

with the external potential

$$V_{\text{ext}}(\boldsymbol{r}) = \frac{1}{2} \left(m\omega_x^2 x^2 + m\omega_y^2 y^2 + m\omega_z^2 z^2 \right) + sE_R \cos^2(kz).$$
(3.27)

Here, the confinement along the radial (x- and y-) direction comes from both the harmonic trap and the optical lattice beams. Changing the depth s of the lattice therefore also changes the radial confinement, but leaves the axial (z-) confinement unchanged.

If the lattice is sufficiently deep, the depth of the lattice wells is larger than the chemical potential and the BEC is split into an array of pancake-shaped BECs, which are coupled to each other due to tunneling between the wells. In this case, the system is best described in a tight-binding picture, where the total wave function of the system is written as a sum of localized wave functions $\Psi_j(\mathbf{r}, N_j)$ centered at the individual lattice sites j,

$$\Psi(\boldsymbol{r},t) = \sum_{j} c_{j}(t) \Psi_{j}(\boldsymbol{r},N_{j}), \qquad (3.28)$$

and $N_j = |c_j|^2$ is the number of atoms at lattice site j. In the case where the chemical potential is much lower than the vibrational level spacing $\hbar \omega_{\text{lat}}$ of the lattice wells, only the lowest level is populated. Excitations along the tightly confined direction are "frozen out" and the dynamics is restricted to the radial, weakly confining direction. We can then write

the localized wave function as a product of a Wannier function along the lattice axis *z* and a density-dependent wave function along the radial axis,

$$\Psi_j(\boldsymbol{r}, N_j) = w(z - jd)\Phi(x, y, N_j). \tag{3.29}$$

Inserting this ansatz into the Gross-Pitaevskii equation and integrating over the spatial coordinates, one arrives at a discrete nonlinear Schrödinger-like equation (DNL) [Sme03],

$$i\hbar \frac{\partial c_j}{\partial t} = J(c_{j-1} + c_{j+1}) + \mu_j^{loc}(c_j)c_j + V_jc_j.$$
(3.30)

In the derivation of this equation, two further important approximations have been made. First, the overlap between wave functions more than one lattice site apart is assumed to be negligible. Second, the local wave functions $\Psi_j(r, N_j)$ only depend implicitly on time through N_j . The radial wave function at a lattice site is assumed to always adapt to the ground state corresponding to the number of atoms in that well, which means that any radial excitations are not taken into account.

The right-hand side of equation (3.30) consists of three terms: The first term describes hopping between neighboring lattice sites. The tunneling matrix element J is calculated as

$$J = -\int \left(\frac{\hbar^2}{2m} \nabla \Psi_j \cdot \nabla \Psi_{j+1} + \Psi_j V_{\text{ext}} \Psi_{j+1}\right) d\mathbf{r}.$$
(3.31)

In principle, this integral depends implicitly on the respective populations N_j at the lattice sites. However, the dependence of J on N_j is very weak [Sme03] and one can use the approximation $\Psi_j(\mathbf{r}, N_j) \approx \Psi_j(\mathbf{r}, N_0)$, where N_0 is the average number of atoms per site. Usually, one does not have to compute the integral directly; for the case of a deep lattice it can be shown that $J = 4\delta$, where δ is the width of the lowest Bloch band [Jak99]. The band structure is easily calculated as outlined in section 3.1.2. For lattices where $s \gg 1$, the tunneling energy can also be approximated by the analytical expression [Blo08]

$$J/E_R = \frac{4}{\sqrt{\pi}} s^{3/4} e^{-2\sqrt{s}}.$$
(3.32)

This approximation should be used with some care – it overestimates J by 30 % for 5 E_R lattice depth and by 15 % for 15 E_R lattice depth. The value of J is plotted in figure 3.5 for various lattice depths.

The second, nonlinear term in equation (3.30) describes the on-site interaction, characterized by the local chemical potential (cf. equations (2.18) and (2.19))

$$\mu_j^{loc} = \int \left(\frac{\hbar^2}{2m} |\nabla \Psi_j|^2 + V_{\text{ext}}(\boldsymbol{r}) |\Psi_j|^2 + g |\Psi_j|^4\right) d\boldsymbol{r}.$$
(3.33)

For an effectively two-dimensional BEC with the wave function (3.29) in a harmonic trap, the chemical potential is, in the Thomas-Fermi approximation [Sme03],

$$\mu_j^{loc} = U_1|c_j| = \sqrt{\frac{m\bar{\omega}_\perp \tilde{g}}{\pi}}|c_j|, \qquad (3.34)$$

where $\omega_{\perp} = \sqrt{\omega_x \omega_y}$ is the effective radial trapping frequency and \tilde{g} is an effective 2D interaction coupling constant given by

$$\tilde{g} = g \int |w(z)|^4 dz \approx \frac{g}{\sqrt{2\pi}\sigma_{\text{lat}}},$$
(3.35)



Figure 3.5: The tunneling matrix element *J* as a function of lattice depth. The solid line is computed from the width of the lowest band in a band structure calculation. The dashed line is the approximation (3.32).

where the Gaussian approximation to the Wannier function has been used. Note that since $\sigma_{\text{lat}} \propto s^{-1/4}$, \tilde{g} and therefore μ_j^{loc} scale weakly with the lattice depth.

The third term describes the external confinement, which gives rise to a potential energy difference between lattice sites. Here this is simply the harmonic trap,

$$V_j = \frac{1}{2}m\omega_z^2 d^2 j^2.$$
 (3.36)

3.1.5 Ground state of a BEC in a lattice

What is the ground state of a BEC in a 1D lattice and a harmonic confinement? One can obtain a stationary solution to the DNL by setting $c_j(t) = g_j e^{-i\mu t/\hbar}$, which leads to

$$\mu g_j = J \left(g_{j-1} + g_{j+1} \right) + \left(\mu_j^{loc}(g_j) + V_j \right) g_j.$$
(3.37)

An analytic solution can be found using the Thomas-Fermi approximation and neglecting the kinetic energy term, which here corresponds to the tunneling term. This leads to the equation

$$\mu = \mu_j^{loc}(g_j) - V_j. \tag{3.38}$$

Using $\mu_j^{loc} = U_1 |g_j|$, the number of atoms at lattice site j is then

$$N_{j} = \begin{cases} \left(\frac{\mu - V_{j}}{g}\right)^{2} = N_{0} \left(1 - \frac{j^{2}}{j_{TF^{2}}}\right)^{2} & \text{where } \mu > V_{j}, \\ 0 & \text{otherwise} \end{cases}$$
(3.39)

When many lattice sites are occupied, the normalization condition $\sum_j N_j = N$ can be replaced by an integral. In a similar manner to the Gross-Pitaevskii equation as discussed in section 2.1.2, several useful relations can now be calculated. One finds [Sme03] that the global chemical potential

$$\mu = \left(\frac{15NU_1^2 \sqrt{m\omega_z^2 d^2/2}}{16}\right)^{2/5}.$$
(3.40)



Figure 3.6: The solid line shows the dependence of the effective scattering length a^* on the lattice depth. The dashed line indicates the harmonic oscillator ground state approximation, equation (3.44). The dash-dotted line has been calculated using the improved Gaussian approximation described by equation (3.19).

This somewhat complicated expression can be rewritten in a familiar form,

$$\mu = \frac{1}{2} \hbar \bar{\omega} \left(\frac{15Na^*}{\sigma_{ho}} \right)^{2/5},\tag{3.41}$$

where $\bar{\omega} = (\omega_x \omega_y \omega_z)^{1/3}$ is the geometrical average of the trapping frequencies for the harmonic confinement, $\sigma_{ho} = \sqrt{\hbar/(m\bar{\omega})}$ the associated harmonic oscillator length and a^* an effective scattering length,

$$a^* = ad \int |w(x)|^4 dx.$$
 (3.42)

Equation (3.41) is equal to the expression for the chemical potential in harmonic trap without a lattice, equation (2.26), except that the scattering length a has been replaced with a^* . This can be understood in the following manner: The lattice has the effect of increasing the local density at the lattice sites. Compared to a wave function with constant density over one lattice site, the interaction energy corresponding to a Wannier function is increased by a factor

$$\frac{E_{\text{int,wannier}}}{E_{\text{int,const.}}} = \frac{g \int |w(x)|^4 dx}{g \int |\Psi_{\text{const.}}|^4 dx} = d \int |w(x)|^4 dx.$$
(3.43)

Thus, on a macroscopic scale, the periodic lattice potential results in an effective enhancement of the interaction energy. If one introduces a "smoothed" macroscopic density, defined as the average density over one lattice well, the usual expressions for chemical potential, density profile and Thomas-Fermi radii described in section 2.1.2 can be used if one replaces a with a^* . A more rigorous discussion of this concept can be found in [Ped01, Kra02].

The dependence of the effective scattering length on the lattice depth is plotted in figure 3.6. For deep lattices, where the Wannier functions approach the harmonic oscillator ground state (equation (3.17)), one can use the analytical result

$$\frac{a^*}{a} = \frac{s^{1/4}}{\pi}.$$
(3.44)

For 10 E_R lattice depth, this approximation overestimates a^* by about 10 %.

It is interesting to consider the effect of adding lattices along the other spatial dimensions. In a 3D lattice, the integral in equation (3.42) has to be carried out over three dimensions instead of one, and thus $a_{3D}^* = (a^*)^3$. A 3D lattice therefore strongly enhances the effective interactions. For a 10 E_R deep lattice, $a^*/a = 2.0$ and $a_{3D}^*/a = 8.4$.

3.2 Technical setup

3.2.1 Lattice setup

A commercial non-planar ring oscillator laser (Innolight Mephisto) provides 2W of 1064 nm light with the narrow bandwidth required to create an optical lattice. This light is split up and distributed to several labs. We have about 700 mW at our disposal, which is used to seed a home-built fiber amplifier using an Ytterbium-doped large-mode-area fiber [Lie03]. When the fiber is pumped with 33 W of 980 nm light from a high-power laser diode (Dilas MIF 980), the seed light is amplified to typically 15 W without loss of spectral quality.

If the amount of seed light coupled into the fiber becomes too small, due to vibrations or accidental blocking of the seed beam, not all energy stored in the fiber is removed. Within milliseconds, the fiber amplifier becomes unstable and will start removing its energy in pulses, which quickly destroys the fiber end. To protect against this, the outgoing laser intensity is monitored with a photodiode behind one of the dielectric mirrors in the beam path. If the intensity changes too quickly, an electronic circuit automatically turns off the pump light.

The 15 W laser light from the fiber amplifier is split up into five beam paths, two beams to create the dimple trap (see section 2.3.4) and three lattice beams to create a 3D lattice. For the experiments described in this work, only one lattice beam is used, creating a 1D lattice.

Figure 3.7 shows the optical setup for the lattice, described in detail in [RK07]. The lattice light is guided to the experiment using a single-mode polarization maintaining fiber, with a Faraday isolator installed directly after the fiber providing protection from the retro-reflected light. A photodiode measures the light intensity after the Faraday isolator. This signal is sent to a servo circuit which controls the intensity using an AOM installed before the fiber. The light is focused onto the atomic sample using a pair of lenses. The beam waist at its focus is approximately 300 μ m. After the focus, the beam is collimated using another lens pair and then retro-reflected by a mirror to create a standing wave. Before the mirror, the vertically propagating dimple beam is overlapped with the lattice beam on a polarizing beam splitter. According to equation (3.2), the lattice depth is

$$V_0 = 0.11 E_R / (I/\text{mW}).$$
 (3.45)

The fairly large beam waist ensures a homogenous lattice depth. For a BEC with Thomas-Fermi radius 30 μ m in the center of the lattice, the lattice depth changes less than 1 % over the sample.

The lattice is often combined with the dimple trap to increase the harmonic confinement. Additionally, the magnetic levitation causes a horizontal antitrapping which slightly weakens the confinement. The potential of the combined traps is

$$V(\mathbf{r}) = \frac{1}{2} \left(m\omega_x^2 x^2 + m\omega_y^2 y^2 + m\omega_z^2 z^2 \right) + V_0 \cos^2(kz),$$
(3.46)

57



Figure 3.7: Illustration of the setup for the optical lattice.

with the trap frequencies in *x*-,*y*- and *z*-direction being

$$\omega_x^2 = \omega_{Dh}^2 + \omega_{Dv}^2 + \omega_{lat,r}^2 - \omega_{lev}^2, \qquad (3.47)$$

$$\omega_y^2 = \omega_{Dv}^2 + \omega_{lat,r}^2 - \omega_{lev}^2, \qquad (3.48)$$

$$\omega_z^2 = \omega_{Dh}^2. \tag{3.49}$$

Here, ω_{Dh} and ω_{Dv} are the trap frequencies of the horizontal and vertical dimple beams, respectively. The antitrapping frequency from magnetic levitation is usually $\omega_{lev} \approx 2\pi \cdot 3$ Hz and the radial confinement of the lattice $\omega_{lat,r} = 2\pi \cdot 2.55 \sqrt{V_0/E_R}$ Hz. The experiments in this thesis are performed with an 8 E_R deep lattice, corresponding to $2\pi \cdot 7.2$ Hz confinement.

Rough alignment of the lattice is done by shining the lattice beam on the cold and spatially large atom cloud produced after Raman sideband cooling. The cloud is sufficiently cold that the dipole trapping effects of the beam can be seen. Fine alignment is done by observing the effect of the lattice beam on the condensate. The beam is adiabatically ramped up and the displacement of the BEC is monitored and minimized, which assures alignment in the imaging plane. Alignment along the imaging axis can also be probed by suddenly switching off an adiabatically raised lattice beam and observing oscillations in the dimple trap. If there are no oscillations, alignment along the imaging axis is assured. Finally, the retroreflected beam is adjusted by optimizing its coupling back into the fiber (the protecting optical diode is slightly de-adjusted such that it allows a small amount of the returning beam to come through).

3.2.2 Lattice depth calibration

The depth of the optical lattice could be calculated from the power and size of the lattice beams. These quantities are difficult to measure precisely and the beam alignment may be imperfect, therefore a way of measuring the lattice depth is needed. We use a method outlined by Denschlag et al. [Den02].

A BEC with momentum q in the lab frame can to a good approximation be described as a plane wave $|q\rangle$. Suddenly switching on the lattice projects the plane wave onto Bloch states,

$$|\Psi(t=0)\rangle = \sum_{n} |\phi_q^{(n)}\rangle \langle \phi_q^{(n)}|q\rangle$$
(3.50)

$$= \sum_{n} c_{n,0}^{*}(q) |\phi_{q}^{(n)}\rangle, \qquad (3.51)$$

according to equation (3.11). While in the lattice, the BEC wave function evolves in time according to

$$|\Psi(t)\rangle = \sum_{n} c_{n,0}^{*}(q) e^{-i\frac{E_{n}(q)}{\hbar}t} |\phi_{q}^{(n)}\rangle.$$
(3.52)

After a time t_h the lattice is suddenly switched off and the Bloch states are projected back onto plane waves. The final state is now a superposition of momentum states separated by $2\hbar k_r$

$$|\Psi(t)\rangle = \sum_{n,l} c_{n,0}^*(q) c_{n,l}(q) e^{-i\frac{E_n(q)}{\hbar}t_h} |q+l\cdot 2\hbar k\rangle, \qquad (3.53)$$

59

which can be observed as diffraction peaks in a time-of-flight measurement. The population of diffraction peak l is

$$p_l(t) = \left| \langle q + l \cdot 2\hbar k | \Psi(t) \rangle \right|^2 = \left| \sum_n c_{n,0}^*(q) c_{n,l}(q) e^{-i\frac{E_n(q)}{\hbar}t_h} \right|^2.$$
(3.54)

The interference between the differently evolving phases leads to oscillations in the diffraction peak populations as a function of t_h .

When starting with a BEC at rest (q = 0), no odd Bloch bands get populated because the corresponding Bloch wave functions are antisymmetric, while the initial wave function is symmetric. Furthermore, for the lattice depths we use in the lab, only bands 0 and 2 are significantly populated². Therefore the population of the diffraction peaks will oscillate with a frequency corresponding to the energy difference between band 0 and 2,

$$f_{2-0} = \frac{E_2 - E_0}{h},\tag{3.55}$$

which is dependent on the lattice depth. This dependence can be computed from a band structure calculation and is plotted in figure 3.8.



Figure 3.8: Left: Energy difference between Bloch band 0 and 2 for quasimomentum q=0. **Right:** A typical lattice depth measurement showing oscillations in population between the $0\hbar k$ and $\pm 2\hbar k$ momentum states. For this particular case, the oscillation frequency is 8.82(3) kHz, corresponding to 7.93(4) E_R lattice depth.

A typical lattice depth calibration is carried out in the following manner: We produce a BEC, turn off the confining dipole trap and tune the scattering length close to zero by switching the magnetic offset field to 17 G. The lattice is then pulsed on for a time t_H and after typically 50 ms of expansion the atom cloud is imaged to reveal the momentum distribution. The BEC is levitated during the whole procedure.

The momentum distribution shows two diffraction peaks at $\pm 2\hbar k$. We count the fraction of atoms in these two peaks and plot against the pulse length. The result is shown in figure 3.8. Note that the amplitude of the oscillation signal decreases significantly slower than in [Den02]. The authors attribute the decay to inhomogeneity of the lattice beams leading to

²The population in band 4 is less than 0.3 % for a lattice depth of 10 E_R . To populate band 4 with more than 5 % requires a depth of 24 E_R .

slightly different oscillation periods over the sample. In our setup the lattice depth is more homogenous due to the large size of our lattice beams, enabling the observation of many oscillation periods and therefore precise lattice depth measurements. 4

Bloch oscillations and interaction-induced dephasing

In the early days of quantum mechanics, the study of electrical conductivity in a crystal lattice led Felix Bloch to predict that when a force acts on a particle in a periodic potential, it will undergo an oscillating motion [Blo29]. In practice however, an electron in a conventional solid will not Bloch oscillate because it will scatter on another electron or a crystal defect before it has time complete to a single oscillation cycle. It was not until 1993 that Bloch oscillations could be observed using semiconductor superlattices, where the larger lattice spacing leads to a much shorter Bloch period [Was93].

Ultracold dilute gases in optical lattices provide an interesting alternative for observing Bloch oscillations, with easy control over many parameters and the ability to directly image the momentum distribution of the atomic sample. Bloch oscillations were first seen with velocity-selected thermal atoms [BD96] and later observed with a BEC [And98, Cri02]. Exactly as for electrons in a solid, particle-particle interactions lead to collisional dephasing and severely limit the number of Bloch oscillations one can observe. One way around this problem is to use a Fermi gas of indistinguishable atoms, where s-wave scattering is prohibited due to the symmetry requirements of the scattering wave function. In this way, about 200 oscillations were observed [Roa04]. With bosons, 4000 oscillations where observed with a thermal cloud of ⁸⁹Sr atoms at low density [Fer06], using the fact that ⁸⁹Sr has a background scattering length that is zero to within a few a_0 [Mic05, Esc08].

The ability to tune interactions using a wide Feshbach resonance gives us an excellent tool to study the interaction-induced dephasing of Bloch oscillations. We have investigated the rate of dephasing of a Bloch oscillating BEC as a function of scattering length. We are able to resolve changes in the scattering length as small as $0.1 a_0$, and when tuning the scattering length to zero, we find that we can follow more than 20000 oscillations [Gus08]. Similar experiments using ³⁹K have been performed at LENS, Italy [Fat08a].

The suppression of atom-atom interactions has potential implications for atom interferometry. A BEC has a very narrow momentum spread, and its small size makes it possible to investigate effects on a μ m scale. However, the large density leads to large interaction effects which limit coherence time and accuracy [Gup02]. One possible way around this limitation is to work with low densities using low atom numbers and weak traps [Ree05], but this 4. Bloch oscillations and interaction-induced dephasing

limits the signal-to-noise ratio. Our technique, instead, allows the direct suppression of the effect of interactions. The large number of Bloch oscillations we could observe allowed us to measure the force acting on the BEC with a sensitivity better than 10^{-6} . This shows that precision measurements with BECs are feasible.

4.1 Theory of Bloch oscillations

When a matter wave in a lattice potential is subject to an external force, it will not be accelerated towards infinity. Instead, the force induces an oscillatory motion, the Bloch oscillation [Blo29]. This can be understood in a band structure picture. A wavepacket in a lattice exposed to a force *F* will gain in quasimomentum as q = Ft. When it reaches the end of the Brillouin zone, because of the periodicity of the band structure (a state with quasimomentum $q + 2\hbar k$ is equivalent to a state with quasimomentum *q*) it will "wrap around" and emerge at the other side of the Brillouin zone. This leads to a periodic motion in quasimomentum space with a period corresponding to the time it takes for the wavepacket to "scan" the whole Brilluoin zone,

$$T_B = \frac{2\hbar k}{F}.\tag{4.1}$$

The oscillation in quasimomentum space also leads to an oscillation in position space with the amplitude [Har04]

$$l_B = \frac{\Delta}{2F},\tag{4.2}$$

where Δ is the width of the lowest band. In experiments with ultracold atoms, this amplitude is usually very small (in the experiments in this chapter, the amplitude is 50 nm) and Bloch oscillations are observed by imaging in momentum space.

When the Bloch oscillating wavepacket reaches the edge of the Brillouin zone, there is a probability of Landau-Zener tunneling into the next higher Bloch band, which effectively leads to the particle being lost from the lattice. The tunneling probability depends on the magnitude of the applied force and is [Mor06]

$$r = e^{-F_c/F}$$
, where $F_c = \frac{V_0^2}{E_R d} \cdot \frac{\pi^2}{32}$ (4.3)

and *d* is the distance between the lattice sites. The experiments described in this thesis are all carried out at a lattice depth $V_0 = 7.9 E_R$. In this case, the critical force $F_c \approx 3 \cdot 10^{-23}$ N. The force acting on the atoms is the gravitational force mg, which is about 15 times smaller than the critical force, which gives $r \approx 4 \cdot 10^{-7}$. Therefore, tunneling into higher bands can be neglected for the rest of the discussion.

Another instructive way of looking at Bloch oscillations is a quantum mechanical treatment in terms of the phase evolution at each lattice site. The Hamiltonian of the system is

$$\hat{H} = \frac{\hat{p}}{2m} + V(x) + Fx,$$
(4.4)

where V(x) is the lattice potential with period *d*. The eigenfunctions to this Hamiltonian are the Wannier-Stark functions¹, which can be constructed as a superposition of Wannier

¹Strictly speaking, the Wannier-Stark states are only eigenstates when tunneling into higher Bloch bands is neglected. See [Har04] and references therein for a more complete discussion.

4.2. Experimental realization

functions [Har04]:

$$\Psi_{\text{WS},j} = \sum_{l} B_{k-j} \left(\frac{Fd}{J}\right) w_k(x), \qquad (4.5)$$

where B_l is the l^{th} Bessel function, w_l is the Wannier function centered over the l^{th} lattice site and J is the tunneling matrix element. In the case of a large force such that $Fd/J \gg 1$, all Bessel functions except B_0 are approximately zero and the Wannier-Stark functions are reduced to Wannier functions. Let us in the following discussion for simplicity assume that we have such a large force, meaning that the Wannier functions are eigenfunctions to the Hamiltonian (a similar argument as below can also be made in the general case).

The total wave function of the system can be written as a sum of Wannier functions localized to the individual lattice sites. In momentum space, the wave function takes the form

$$\Phi(p,t) = \sum_{j} c_{j} w_{j}(p) = w_{0}(p) \sum_{j} c_{j} e^{-ipjd/\hbar},$$
(4.6)

where the relation $w_j(p) = w_0(p)e^{-ipjd/\hbar}$ comes as a result of the translation $w_j(x - jd) = w_0(x)$. For a non-interacting system, the phase on each lattice site evolves according to the site's potential energy,

$$\hbar \frac{d\phi_j}{dt} = Fjd. \tag{4.7}$$

The total wave function can then be written as

$$\Phi(p,t) = w_0(p) \sum_{j} c_j(0) e^{iFdjt/\hbar} e^{-ipjd/\hbar}$$
(4.8)

$$= w_0(p) \sum_{j=1}^{n} c_j(0) e^{-i(p-Ft)jd/\hbar}.$$
(4.9)

As discussed in section 3.1.3, the wave function in quasimomentum space is found by dividing by the Wannier function (in this case particularly simple), which leads to

$$\Phi(q,t) = \sum_{j} c_j(0) e^{-i(q-Ft)jd/\hbar}.$$
(4.10)

It is now clear that an initial quasimomentum distribution f(q) will after a time t have evolved to f(q - Ft), e.g applying a force results in a translation of the wave function in quasimomentum space.

4.2 Experimental realization

4.2.1 Observation of Bloch oscillations

Our protocol for observing Bloch oscillations is illustrated in figure 4.1. We start with a nearly pure condensate with typically $1 \cdot 10^5$ atoms in the dimple trap with trap frequencies $2\pi \cdot (11, 6, 10)$ Hz. The BEC has Thomas-Fermi-radii (13,33,12) µm and a peak density of about 10^{13} cm⁻³. A 1D optical lattice oriented along the vertical axis is then ramped to 7.9 E_R lattice depth exponentially in 1000 ms, and at the same time the power in the vertically propagating dimple trap beam is increased. The confinement is then increased to $2\pi \cdot (15, 12, 10)$ Hz. The

4. Bloch oscillations and interaction-induced dephasing



Figure 4.1: The experimental protocol used to observe Bloch oscillations. The BEC is adiabatically loaded into the optical lattice. Bloch oscillations are then induced by turning off the magnetic levitation and the vertical confinement. After a hold time t_H , the atomic sample is adiabatically released from the lattice, followed by ballistic expansion with the scattering length set to zero.

slow ramp is adiabatic not only with respect to the tight lattice confinement, but also with respect to the radial trapping. This ensures that the condensate does not get excited along the horizontal direction. About 50 lattice sites are populated.

Next, the magnetic levitation and the vertical confinement is switched off. Our relatively large lattice beams produce a very weak horizontal confinement on their own, 8.7 Hz for 7.9 E_R lattice depth. We therefore do not turn off the vertical dimple beam, which copropagates with the lattice. With magnetic levitation turned off, the condensate starts to Bloch oscillate due to the gravitational force. It is left to evolve in the lattice for a hold time t_H , after which the levitation is switched on again.

How can we now measure the quasi-momentum distribution? If the lattice is ramped down on a time scale that is slow compared to the lattice trap frequency, but fast enough that the quasimomentum distribution does not change, a state with quasimomentum q is adiabatically mapped to a state with free particle momentum p = q [Kas95, Gre01]. The free particle momentum can then be measured in ballistic expansion. We ramp down the lattice and the vertical dimple beam in 300 µs and take an absorption picture after a period of free expansion, typically 80 ms. The scattering length is switched to zero when the trap is turned off. This enables us to directly image the quasimomentum distribution without any broadening due to interaction energy being converted into kinetic energy, which otherwise would greatly reduce the contrast of the Bloch oscillations [Roa04].

The result of this experiment is shown in figure 4.2. The quasimomentum increases with increasing hold time t_H until it reaches lower edge of the first Brillouin zone. There it "wraps around" and appears at the other edge. Note that some parts of the quasimomentum distri-



Figure 4.2: A sequence of absorption images revealing the evolution of the quasimomentum distribution during a Bloch oscillation. Every image is the average over 7 snapshots and the sequence shows the evolution over 575 µs. The dashed lines indicate the extent of the first

bution appears to be outside of the Brillouin zone. This is due to the fact that the lattice cannot be ramped down perfectly adiabatically with the respect to the vibrational trap frequency, since the trap frequency approaches zero at the end of the ramp. The mapping of quasimomentum to momentum is then not perfect.

4.2.2 Momentum broadening due to interactions

Brillouin zone.

If interactions are negligible during Bloch oscillations, the dominant factor influencing the phase evolution of the wave function at a lattice site is the potential due to the external force. In this case, Bloch oscillations would go on forever. When interactions are present, the phase at each lattice site will in addition be influenced by the local interaction energy, which is different from site to site. This gives rise to dephasing, which causes a broadening of the quasimomentum wave function and limits the number of oscillations one can observe.

We study this effect by adjusting the scattering length during Bloch oscillations in the range² from 0 to 300 a_0 . To keep the initial density constant, the lattice is always loaded at a scattering length of 210 a_0 . As a quantitative measure of the dephasing, we measure the rmsdiameter Δp of the momentum distribution at instants in time when the average momentum is zero, i.e. after an integer number of oscillations N. Figure 4.3a shows how Δp evolves with time for different scattering lengths. When the scattering length is switched to $\approx 0 a_0$, no broadening is observed within the first 300 oscillations. For larger scattering lengths Δp increases linearly with time until a saturation value of about 1.3 $\hbar k$, which corresponds to the first Brillouin zone being completely filled.

Figure 4.3b shows how Δp depends on the scattering length for a fixed number of oscillations. It appears that Δp increases with the square root of the scattering length. An initially linear increase with both time and interaction strength has been predicted in a pure 1D model [Wit05]. As discussed in section 3.1.4, our 3D system can be described with an effective 1D model where the effective interaction, characterized by the local chemical potential μ^{loc} defined in equation (3.34), scales with the square root of the scattering length.

²Note that it is not possible to switch to a significant negative scattering length, since the BEC becomes unstable and undergoes a "Bose-Nova" style implosion [Don01] for scattering lengths $\leq -5 a_0$.

4. Bloch oscillations and interaction-induced dephasing



Figure 4.3: (a) Evolution of the rms-diameter Δp of the momentum distribution for different values of the scattering length. a = 0, 25, 50, 100 and $300 a_0$ from bottom (filled circles) to top (open squares). (b) Δp as function of scattering length after a fixed number of Bloch oscillations N = 1 (filled circles), 25 (filled squares), 50 (filled diamonds), 100 (open circles), 150 (open squares) and 200 (open diamonds). The solid lines are numerical solutions to equation (3.30), see text.

Hence our observations are in agreement with the predictions in [Wit05].

We have also simulated the system by solving the DNL (equation (3.30)) numerically. The result is indicated by the solid lines in figure 4.3 and is in good agreement with the data. There is a systematic discrepancy for the N = 50 data in figure 4.3b, where the measured dephasing is stronger than in the simulation. We attribute this to horizontal dynamics not included in the model. When the Bloch oscillations are started, the scattering length is suddenly switched to a new value and the horizontal trapping beam is switched off, causing the trap frequencies to change. Both of these changes excite a horizontal breathing mode, which causes the density to be modulated in time. This causes a modulation of the effective interaction, which results in a change in the rate of dephasing. More details about the simulation can be found in [Mar07a].

One can learn more about the dephasing by examining the shape of the quasimomentum distribution. When the BEC wave function completely fills the first Brillouin zone, interference fringes appear. A detailed investigation of this phenomenon and its implications will be the subject of chapter 5.

4.2.3 Precise determination of the scattering length zero crossing

If the atoms are left to Bloch oscillate in the lattice for a long time, even a very small amount of interaction causes dephasing and broadens the momentum distribution of the BEC. This can be used as a sensitive tool to determine at which magnetic field strength the scattering length is zero.

We seek the scattering length zero crossing by measuring the width of the quasimomentum distribution after the BEC has been Bloch oscillating for 4 s. To be more sensitive to small interaction strengths, the density of the BEC was increased to about $7 \cdot 10^{13}$ cm⁻³ for this measurement.

The result of this experiment is shown in figure 4.4. There is only a small range of ≈ 20 mG centered around 17.119 mG where the BEC wave function is not spread out over the whole Brillouin zone. Within this range, even a change in the scattering length as small as 0.1 a_0 has an observable effect on the BEC. Since we can also reach scattering lengths up

4.2. Experimental realization



Figure 4.4: Width of the quasimomentum distribution after 6951 Bloch oscillations (4 s evolution time) in a compressed trap at different values of the magnetic field. The solid line is a Gaussian fit to the data, to guide the eye and to determine the position of the dephasing minimum. The minimum is found at 17.119(2) mG. Above the plot is a scale indicating the scattering length. This scale has been chosen such that the zero agrees with the value 17.13 mG inferred from this measurement, including a correction due to magnetic dipole-dipole interactions, and the slope is 61 G/ a_0 , inferred from coupled-channel scattering calculations performed at NIST [Jul04].

to 2000 a_0 , this shows that we can reproducibly change the scattering length over more than four orders of magnitude.

Although the point of minimum dephasing can be inferred with high precision from the data in figure 4.4, this does not automatically translate into an accurate and precise measurement of the scattering length zero crossing. I will now discuss the statistical error and the main factor limiting accuracy, magnetic dipole-dipole interaction.

Magnetic field measurement and statistical error

The magnetic field value was measured by microwave spectroscopy. A 200 ms long microwave pulse was applied to a cold thermal gas in the same trap configuration that was used to measure the dephasing minimum. If the pulse is resonant with the ($|F = 3, m_F = 3\rangle \rightarrow |F = 4, m_F = 4\rangle$) transition, heating due to spin relaxation leads to atom loss. The magnetic field can be determined from the resonance frequency using equation (2.37), as discussed in section 2.2.8. The long pulse duration was used to make sure that any short-time fluctuations in the magnetic field are averaged out.

The magnetic field was measured for two different control voltages given to the Helmholtz coil control circuit. The field *B* corresponding to a control voltage V is calculated from a simple linear interpolation

$$B(V, B_2, B_1) = \frac{B_2 - B_1}{V_2 - V_1} (V - V_2) + B_2,$$
(4.11)

where B_1, B_2 is the field measured for the control voltages V_1, V_2 . The total statistical error in determining *B* is then

$$\Delta B(V_0, B_1, B_2) = \sqrt{\left(\frac{\partial B}{\partial V_0} \Delta V_0\right)^2 + \left(\frac{\partial B}{\partial B_1} \Delta B_1\right)^2 + \left(\frac{\partial B}{\partial B_2} \Delta B_2\right)^2}.$$
 (4.12)

4. Bloch oscillations and interaction-induced dephasing

We measure B_1 =17.1426(16) G for V_1 = 1.8 V and B_2 = 17.0937(12) G for V_2 = 2.2 V. A Gaussian fit to the data in figure 4.4 gives that the minimum dephasing is at a control voltage V_0 = 1.992(36) V. From this we can calculate that dephasing is minimized at a magnetic field of 17.119(2) G. The uncertainty from the fitting process contributes about as much to the final error as the uncertainty from the magnetic field measurement.

Shift due to magnetic dipole-dipole interaction

Apart from the contact interaction characterized by the scattering length, there are other forms of atom-atom interaction. The strongest of these is the magnetic dipole-dipole interaction [Men08] which is described by a potential

$$V_{dd}(\mathbf{r}) = -\frac{\mu_0 \mu^2}{4\pi} \left(\frac{3\cos^2 \theta - 1}{r^3}\right),$$
(4.13)

where μ_0 is the Bohr magneton, μ is the magnetic moment of the atom, r the distance between the two interacting dipoles and θ the angle between the direction joining the two dipoles and the dipole orientation (assuming that the dipoles are aligned in the same direction). The relative strength of the dipolar interaction compared to the contact interaction is characterized by the dimensionless parameter

$$\epsilon_{dd} = \frac{\mu_0 \mu^2 m}{12\pi\hbar^2 a}.\tag{4.14}$$

In most cases $\epsilon_{dd} \ll 1$, for example Cs with $a = 210 a_0$ yields $\epsilon_{dd} \approx 0.003$, and the dipolar interaction is completely negligible. When we tune the scattering length close to zero, this is no longer the case. For $a = 0.5 a_0$, $\epsilon_{dd} \approx 1.2$. Therefore, the total on-site interaction is not zero when the scattering length is zero, but rather where the contact and dipolar interactions cancel each other out. The scattering length \bar{a} where this happens has been calculated in [Fat08b], approximating the BEC wave function at a lattice site with a Gaussian with radius σ_{\perp} and σ_{lat} in the radial and axial lattice directions, respectively. For the case where the atomic dipoles are oriented along the lattice axis as in our system, it is

$$\bar{a} = -\frac{\mu_0 \mu^2 m}{6\pi\hbar^2} \left(1 - \frac{3\pi}{4} \frac{\sigma_\perp}{\sigma_{\text{lat}}} \right)$$
(4.15)

in the limit $\sigma_{\perp} \ll \sigma_{\text{lat}}$. In our experiment, we excite considerable motion in the radial direction when we switch to a new scattering length during the Bloch oscillations. However, the dependence on the radial size of the BEC is very weak for our experimental conditions, and equation (4.15) should therefore still be valid. For our system, $\bar{a} = -1.13 a_0$.

Since the magnetic dipole-dipole interaction is long-range, it does not only change the on-site interaction. It also gives rise to a non-negligible inter-site interaction term, which was calculated numerically in [Fat08b]. Under the assumption that the relative strength of the inter-site and the on-site interaction is the same for our system, we can estimate the scattering length for which the dephasing is minimized to be $\bar{a} \approx -0.7a_0$. This means that the scattering length zero crossing is at 17.13 mG, offset by 11 mG from the point of minimum dephasing.

4.2.4 Limit of vanishing interaction - long lived Bloch oscillations and measurement of local gravity

Minimizing the interaction strength allows us to observe a large number of Bloch oscillations. As shown in figure 4.5, more than 20000 oscillations can be clearly observed. This corresponds to a coherent absorption and emission of 40000 photons over more than 10 seconds and is the largest number of Bloch oscillations that have been observed.



Figure 4.5: Absorption pictures showing long-lived Bloch oscillations. The pictures are taken after 27 ms of free expansion.

To determine the Bloch oscillation period, we fit two Gaussian peaks to the quasimomentum distribution. The position of the strongest of these peaks can be seen as a measure of the phase of the Bloch oscillation. Figure 4.6 shows how the phase evolves with time. A fit to the data with a sawtooth pattern yields an oscillation period of 0.5753807(5) ms – a relative precision of $9 \cdot 10^{-7}$. To rule out any effects due to gradients caused by the switching of the magnetic fields, the first 100 ms of data were excluded from the fit. Towards the end of the plot, the data points show more scatter. We believe that this is not due to shot-to-shot fluctuations in the actual phase, but rather that the broadening of the quasimomentum distribution makes the fitting procedure more unreliable.

Knowing the Bloch period, local gravity can be determined from equation (4.1). The wavelength of the lattice light was measured to 1064.4946(1) nm with a commercial wavemeter (High Finesse WS7), which gives a local gravity $g = 9.803821(9) \text{ m/s}^2$. The high sensitivity reached clearly shows the potential for precision measurements with BECs with tunable scattering length.

Systematic errors

The error given for local gravity is statistical only. The goal of this experiment was to show that it is possible to make precision measurements with a BEC with tunable scattering length, not to make an accurate measurement of local gravity, and we have made no effort to evaluate the systematic errors. We expect that the main factor limiting accuracy is stray magnetic 4. Bloch oscillations and interaction-induced dephasing



Figure 4.6: The position of the strongest peak in the momentum distribution as a function of time. The solid line is a fit to the data, from which a Bloch period of 0.5753807(5) ms can be extracted.

field gradients, which cause a supplemental force indistinguishable from gravity. In cases where one is content with measuring a *change* in force to high accuracy, this error will cancel out in a differential measurement.

Another cause of systematic error is misalignment between the lattice direction and the force direction. The Bloch oscillation period is inversely proportional to the component of the force acting along the lattice direction, and a 1 mrad misalignment would cause a systematic error of $5 \cdot 10^{-7}$. For the current case where we measure gravity, the beam alignment could be done by slightly tilting the optical table and measuring the force as a function of table tilt.

Limitations

We have shown that when we minimize interactions using a Feshbach resonance, we can observe more than 20000 Bloch oscillations. What limits the number of oscillations we can observe? Although we have not investigated this in detail, I will list some effects that could contribute.

- Magnetic field noise Since the amount of phase picked up at a lattice site due to interactions is proportional to the scattering length (see section 5.1.1), the effect of magnetic field fluctuations around an average value will cancel out with time, as long as the density is constant. This means that magnetic field noise on a time scale faster than any induced density modulations. In our system, there is a density modulation caused by the transverse excitation which is on the time scale of the trap frequencies (≈ 10 Hz). Noise with much higher frequency than this should therefore not influence the dephasing. This means that we are mainly sensitive to long-term drifts in the magnetic field.
- Magnetic dipole-dipole interactions between lattice sites As discussed in chapter 4.2.3, the magnetic dipole-dipole interaction is not negligible when the scattering length is sufficiently small. The on-site dipolar interaction energy has the same dependence of the atom number as the contact interaction and it is possible to have them cancel each other out on every lattice site. However, the dipole-dipole interaction is long-range and there is also an interaction energy due to inter-site interactions [Fat08b]. This term cannot be canceled out and no matter the scattering length, there will always be a small residual interaction leading to dephasing.
4.3. Possibilities for precision measurements with a BEC at zero scattering length

However, there is a way around this limitation. If the atomic dipoles are oriented at an angle $\theta = 54.7^{\circ}$ to the lattice axis, the dipolar interaction energy will average to zero [Gio02]. This could be achieved by an appropriate orientation of the magnetic offset field.

• **Photon scattering** Although the lattice is very far detuned from the atomic resonance, spontaneous emission cannot be neglected on long timescales. Scattering of photons leads to heating and decoherence. It has been shown [Kol02] that in the presence of spontaneous emission, the amplitude of the oscillation of the mean momentum will decay exponentially with a time constant given by the photon scattering rate. We are using a lattice depth of 8 E_R , which at 1064 nm wavelength corresponds to a scattering rate of 0.2 photons/min. This is too slow to explain the decay we are seeing.

Heating due to photon scattering is a fundamental limit to the number of Bloch oscillations one can observe. It is not possible to use a lower lattice depth as this leads to atom losses due to an increased tunneling rate out of the lattice. The way around this limitation is instead to either use a further red-detuned lattice (which will also give a shorter Bloch period due to the smaller lattice k) or to use a lattice detuned to the blue side of the atomic resonance, where the atoms are trapped in scattering minima.

Instead of increasing the number of Bloch oscillations one can observe, the Bloch period precision can be improved by a better measurement of the oscillation phase. We determine the momentum distribution of the BEC by absorption imaging, which is somewhat crude and can reach an estimated resolution of about $\hbar k/200$ if long expansion times are used. Cladé et al. have demonstrated [Cla05, Cla06b] a technique to precisely measure the momentum distribution by employing a velocity-selective Raman transition into another hyperfine state and subsequent spin-dependent detection. A sensitivity of $\hbar k/10000$ was reached with 160 shots.

4.3 Possibilities for precision measurements with a BEC at zero scattering length

The results in the previous section show the potential for doing precision measurements with a BEC with tunable interactions. To further argue this point, I will analyze a scheme for measuring the fine structure constant with a BEC interferometer that could be implemented with the current setup and estimate the main error sources and challenges of such an experiment.

4.3.1 Measuring the fine structure constant with a contrast interferometer

The fine structure constant α is one of the fundamental constants of nature, describing the strength of the electromagnetic interaction in the framework of QED. The CODATA-2006 [Moh08] recommended value of α has an uncertainty of 0.68 ppb, which makes it one of the most precisely measured fundamental constants. However, as illustrated in figure 4.3.1, the recommended value is to a large extent determined by a single experiment, a measurement of the electron magnetic moment carried out at Harvard [Odo06] and an accompanying elaborate QED calculation [Gab06] (0.7 ppb). The danger in having only one method that defines

4. Bloch oscillations and interaction-induced dephasing

the value of a constant is illustrated by the fact that the Harvard 2006 value later had to be adjusted due to a correction of the QED analysis [Gab07]. Recently, the Harvard group published an improved measurement [Han08], reaching an uncertainty of 0.37 ppb.

The best α measurements carried out with other methods come from photon recoil measurements, using Bloch oscillations of thermal Rb atoms [Cla06a, Cla06b] (6.7 ppb) and an atom Ramsey-Bordé interferometer with cold thermal Cs atoms [Wic02, Ger06] (8.0 ppb). These measurements have an uncertainty more than a magnitude larger than the Harvard measurement and there is a clear need for a better alternative α measurement to confirm the Harvard findings. Since the Harvard method needs a QED calculation to arrive at a value for α , an independent α measurement is at the same time a test of QED.



Figure 4.7: Upper: The results of the most recent α measurements. From [Han08]. Lower: The values used to determine the CODATA-2006 value of α . From [Moh08].

I will here briefly describe a method to measure α with a BEC interferometer, as outlined and demonstrated in [Gup03, Gup02]. The method builds on the fact that the fine structure constant can be written as

$$\alpha^2 = \frac{2R_\infty}{c} \frac{h}{m_e} = \frac{2R_\infty}{c} \frac{m_{\rm Cs}}{m_e} \frac{h}{m_{\rm Cs}}.$$
(4.16)

The Rydberg constant R_{∞} is known to better than 7 ppt [Moh08], the Cs mass m_{Cs} has an uncertainty of 0.18 ppb [Wap03, Aud03] and the electron mass m_e is known to 0.42 ppb

4.3. Possibilities for precision measurements with a BEC at zero scattering length

[Moh08]. This means that a precision measurement of \hbar/m_{Cs} would result in a competitive determination of the fine structure constant. \hbar/m_{Cs} can be determined by measuring the photon recoil frequency ω_{rec} , the frequency corresponding to the kinetic energy of a Cs atom recoiling from the absorption of a photon with wave vector k,

$$\hbar\omega_{\rm rec} = \frac{\hbar^2 k^2}{2m_{\rm Cs}}.\tag{4.17}$$

k can be measured very accurately for near-resonant light by referencing it to the Cs D2 line, which has been measured with an uncertainty of 110 kHz or 0.4 ppb [Ude00]. Alternatively, far-detuned light that is referenced to a calibrated frequency comb could be used.



Figure 4.8: (a) 2-way interferometer sensitive to the photon recoil, resulting in a moving density wave. **(b)** Extension to a 3-way contrast interferometer. From [Gup03].

An interferometer scheme for measuring ω_{rec} is illustrated in figure 4.8(a). A standing wave with wave vector k is used to split a BEC into two momentum states, $|0\hbar k\rangle$ and $|2\hbar k\rangle$, using a Bragg $\pi/2$ -pulse. After a time T, a second order Bragg π -pulse is used to reverse the momentum, transferring the $|+2\hbar k\rangle$ state into $|-2\hbar k\rangle$. This causes the two parts to recombine after a time 2T. The two wavepackets will interfere and form a moving density wave . The phase of this wave at time 2T will depend on the relative phase of the wavepackets. The wavepacket in path 2 has been moving with a velocity $2\hbar k$ with respect to path 1, which corresponds to an additional kinetic energy $4\hbar\omega_{\text{rec}}$. Path 2 will then pick up an extra phase

$$\delta \Phi = 8\omega_{\rm rec}T + \phi_1 + 2\phi_2, \tag{4.18}$$

where ϕ_1 and ϕ_2 are the phases of the optical lattice at the time of the two Bragg pulses.

This idea can be extended to a three-way interferometer as shown in figure 4.8(b). The initial splitting of the BEC can be done using a Kapitza-Dirac pulse, creating momentum states $|-2\hbar k\rangle$, $|0\rangle$ and $|2\hbar k\rangle$. When the states are recombined, there will be two density waves moving in opposite direction, which will interfere and create a standing wave. Both the frequency and the phase of this standing wave will depend on the recoil frequency. The standing wave can be probed by scattering a weak light pulse off the density wave, which acts as a matter wave grating. Due to the symmetry of the setup, many error sources will be canceled out. Especially, the standing wave will not be sensitive to the phases ϕ_1 and ϕ_2 and the measurement will therefore not be sensitive to mirror movement between the splitting and recombination pulses. The intensity of the scattered probe pulse, proportional to the

4. Bloch oscillations and interaction-induced dephasing

amplitude of the matter wave grating, will around the time 2T vary as

$$I(t) \sim \cos^2(8N^2\omega_{\rm rec}T + 4N^2\omega_{\rm rec}(t - 2T) + \phi_{\rm off}),$$
 (4.19)

where ϕ_{off} describes a possible phase offset due to for example diffraction phase shifts or light shifts. A measurement of the recoil frequency can then be done by measuring the phase Φ of the density modulation at time t = 2T for varying separation times T, using the relation

$$\Phi = 8\omega_{\rm rec}T + \phi_{\rm off}.\tag{4.20}$$

The scheme can be extended by using higher order Kapitza-Dirac and Bragg pulses to create higher momentum states. The energy (and therefore the measured phase) scales quadratically with the number of photon recoils and the phase ϕ will for N^{th} order pulses simply be

$$\Phi = 8N^2 \omega_{\rm rec} T + \phi_{\rm off}. \tag{4.21}$$

Recently, 12th order Bragg pulses have been demonstrated [Mül08]. The quadratic scaling is an advantage compared to most other interferometer schemes to measure the photon recoil [Wic02, Cla06b], which scale linearly with the number of scattered photons. Other advantages are that the atoms in all interferometer paths are in the same internal state, which makes many systematics common-mode, like for example the AC Stark shift.

4.3.2 Estimates of statistical and systematic errors

The experiment by Gupta et al. [Gup02] was a proof-of-principle experiment, intended to demonstrate the contrast interferometry method. The measured value of α had a statistical error of 7 ppm and deviated by a factor of $2 \cdot 10^{-4}$ from the accepted value. The sensitivity was limited by uncontrollable phase shifts due to the BEC mean-field energy and laser intensity noise, and the accuracy was limited by the mean-field energy [Gup03]. In our setup, we can tune the scattering length to zero and strongly suppress the atom-atom interactions. I will here estimate what errors would affect an α measurement with our system.

Shot noise

Phase fluctuations due to atom shot noise is a fundamental limit to the precision one can obtain in a single shot. Assuming that $N_{\rm at} = 10^5$ atoms contribute to the signal, the phase error is [Gup03] $\Delta \Phi_{\rm sn} \approx 2/\sqrt{N_{\rm at}} \approx 6.3$ mrad. The total phase measured for N = 12 and interferometer time T = 50 ms is $\Phi = 8N^2\omega_{\rm rec}T \approx 7.5 \cdot 10^6$ rad and $\Delta \Phi_{\rm sn}/\Phi \approx 8.5 \cdot 10^{-9}$. To get down to 1 ppb precision, one would need to average over ≈ 70 shots, which would take about 12 minutes in the current setup.

Laser intensity fluctuations

If the intensity of the standing wave changes so that the Bragg pulse is not a perfect π pulse, not only does the population in the different interferometer arms change, there is also a change in the offset phase ϕ_{off} in equation (4.21). Shot-to-shot drifts in laser intensity will then translate into fluctuations of the measured interferometer phase and affect the precision. This error can be estimated to [Büc03b]

$$\Delta \Phi_{\rm int} \approx \frac{4\pi}{3} (1+\delta)^2, \tag{4.22}$$

4.3. Possibilities for precision measurements with a BEC at zero scattering length

where δ is the relative drift in laser intensity. An intensity variation as small as 0.2 % would still change the interferometer phase by 17 mrad, leading to a relative phase error of $2.3 \cdot 10^{-8}$ for N = 12 and T = 50 ms.

Gouy phase

For a plane wave, the wave vector is exactly $k = 2\pi/\lambda$. For a Gaussian beam, the phase along the beam axis evolves as $\phi = kz - \zeta(z)$ where $\zeta(z) = \tan^{-1}(\frac{z}{z_R})$ is the so called *Gouy phase* $(z_R \text{ is the Rayleigh length})$ [Sal91]. This means that the effective wave vector $k_z^{\text{eff}} = \frac{d\phi}{dz}$ is not equal to k. At the beam focus, it is shifted such that

$$\frac{k_z^{\text{eff}} - k}{k} \approx \frac{2}{k^2 w_0^2} \tag{4.23}$$

for a beam with waist (1/ e^2 -radius) w_0 . Since $\omega_{\text{rec}} \propto k^2$, the recoil frequency is thereby shifted by

$$\frac{\Delta\omega_{\rm rec}}{\omega_{\rm rec}} = 2\frac{\Delta k}{k} \approx 1.9 \cdot 10^{-8} \tag{4.24}$$

for a beam with 2 mm waist. One can get a smaller Gouy phase shift by shifting the interaction region away from the focus, but then the wavefront is curved which also leads to a shift of the effective *k* vector. In [Cla06b], the Gouy phase shift and the wavefront curvature for a 2 mm beam lead to a systematic error of 16.4 ppb with an uncertainty of 8 ppb. Although a BEC is much smaller than the $\approx 800 \ \mu m$ diameter thermal cloud in [Cla06b] and the wavefront curvature therefore should be less of a problem, systematics due to the curved wavefront of the Gaussian beam is one of the main hurdles that needs to be overcome to reach an accuracy on the ppb level.

Mean-field shift

The mean-field energy and chemical potential for a condensate was discussed in section 2.1.2. For a cloud with peak density n and scattering length a, the chemical potential for a BEC in the Thomas-Fermi regime released from a harmonic trap is according to equations (2.19), (2.20) and (2.25)

$$\mu = \frac{16\pi\hbar^2 na}{7m}.\tag{4.25}$$

Typical values for a Cs BEC are a density of $n = 10^{13} \text{ cm}^{-3}$ and a scattering length $a = 210 a_0$, which would lead to the phase evolving with frequency $\omega_{\text{int}} = \mu/\hbar \approx 2\pi \cdot 60 \text{ Hz}$. This can be compared with the ¹³³Cs recoil frequency $\omega_{\text{rec}} \approx 2\pi \cdot 2070 \text{ Hz}$, and one immediately sees that the mean-field energy can cause large errors. Any imbalance between the interferometer arms will lead to a shift in the measured interferometer phase. Shot-to-shot atom number fluctuations will limit the precision and a systematic imbalance will lead to a systematic phase shift. If the imbalance is x, the phase shift will be

$$\frac{\Delta\Phi_{\rm int}}{\Phi} = \frac{x\omega_{\rm int}2T}{8N^2\omega_{\rm rec}T} = \frac{x\omega_{\rm int}}{4N^2\omega_{\rm rec}}$$
(4.26)

With an imbalance x = 5%, $\Delta \Phi_{int}/\Phi$ would be $2.5 \cdot 10^{-6}$ for N = 12.

4. Bloch oscillations and interaction-induced dephasing

In the experiment of Gupta et al., the mean-field shift was probably the main factor limiting both precision and accuracy. With the control over the scattering length demonstrated in this chapter, this error can be suppressed by many orders of magnitude. Switching the scattering length to $0.1 a_0$ before the interferometer pulses are applied would reduce the error to $\frac{\Delta \Phi_{\rm mf}}{\Phi} \approx 1 \cdot 10^{-9}$. Reducing the density could lower the error even more.

It is necessary to have a very homogenous magnetic field to reach these low scattering lengths during the whole evolution time of the interferometer. The outer arms of the interferometer are up to $350 \,\mu\text{m}$ apart 50 ms after splitting the BEC and the magnetic field needs to be constant on the mG level over this distance.

Magnetic dipole-dipole interaction

As discussed in section 4.2.3, the magnetic dipole-dipole (MDD) interaction becomes nonnegligible when the scattering length is small. The ratio of mean-field energy from the contact interaction and the mean-field energy from the MDD interaction is [Gio03]

$$\frac{E_{\rm mf,MDD}}{E_{\rm mf,contact}} = \epsilon_{dd} f(\kappa) \left(\frac{3\cos^2\varphi - 1}{2}\right), \tag{4.27}$$

where ϵ_{dd} is the dimensionless parameter characterizing the relative strength of the MDD interaction defined in equation (4.14), κ is the aspect ratio of the BEC and φ is the angle between the magnetic field and the BEC axis. $f(\kappa)$ is a function that is zero for a spherical BEC, approaches 1 in the limit of an elongated cigar-shaped BEC and -2 in the limit of a pancake-shaped BEC. For $0.1 a_0$ scattering length, $\epsilon_{dd} \approx 6$. For a cigar-shaped BEC with aspect ratio 2, $f(\kappa) \approx 0.5$ and the mean-field shift due to MDD interactions is about 3 times larger than the contact interaction shift, or $3 \cdot 10^{-9}$. This could be reduced by either creating a more spherical BEC (in a crossed dipole trap, that would require elliptical beams) or changing the angle of the magnetic field with regard to the BEC axis.

Magnetic field inhomogeneity

For symmetry reasons, errors due to stray magnetic bias fields and gradients are suppressed. The measurement is only sensitive to the magnetic field curvature B''. The energy due to the curvature is $E_B(t) = -\frac{3}{4}\mu_B B'' x(t)^2$ and the extra phase picked up is therefore

$$\Delta\Phi_B = -\frac{2}{\hbar} \int_0^T \frac{3}{4} \mu_B B'' \left(\frac{N\hbar kt}{m}\right)^2 dt$$
(4.28)

$$= -\frac{\mu_B B'' N^2 \hbar k^2 T^3}{2m^2}$$
(4.29)

$$= -\mu_B B'' \frac{N^2 \omega_{\rm rec} T^3}{m},$$
(4.30)

assuming that the velocity change due to the curvature is negligible. This gives a relative error of

$$\frac{\Delta\Phi_B}{\Phi} = \frac{\mu_B B'' T^2}{8m}.\tag{4.31}$$

It should be possible to null the magnetic field curvature to the order of 1 G/cm^2 , which would lead to a relative error of $\approx 1.3 \cdot 10^{-10}$ for an interferometer time *T* of 50 ms. Since the

Error source	Estimated relative error (ppb)
Atom shot noise	8.5
Laser intensity fluctuations	23
Mean-field shift $(210 a_0)$	2500
Mean-field shift $(0.1 a_0)$	1
Magnetic dipole-dipole interaction	3
Gouy phase shift	19
Magnetic field curvature	0.13
Beam alignment	0.9

4.3. Possibilities for precision measurements with a BEC at zero scattering length

 Table 4.1: Summary of the estimated statistical and systematic errors affecting the contrast interferometer.

error scales quadratically with *T*, systematics due to the field curvature might be a limiting factor for longer interferometer times.

It is also worth noting that this systematic prohibits the use of levitation during the measurement. According to equation (2.32), a levitating gradient automatically leads to a curvature of $\approx 7 \cdot 10^8$ G/cm², which would lead to an error of almost 10 % for T = 50 ms! Even for T = 1 ms, the relative error would be more than 10^{-5} .

Beam alignment

The standing wave needed for the Kapitza-Dirac and Bragg pulses is created by two counterpropagating beams. A misalignment of these beams by an angle θ will change the effective wave vector k by

$$\frac{\Delta k}{k} = 1 - \cos\frac{\theta}{2} \approx \frac{\theta^2}{2}.$$
(4.32)

From equation (4.17) we can deduce a relative error

$$\frac{\Delta\omega_{\rm rec}}{\omega_{\rm rec}} = 2\frac{\Delta k}{k} = \theta^2.$$
(4.33)

Other experiments [Cla06a] have demonstrated a misalignment below 30 µrad, which would correspond to a relative error of $9 \cdot 10^{-10}$. If necessary, active beam alignment [Mül05] could further reduce this error.

Summary

The estimated statistical and systematic errors are compared in table 4.3.2. Although this is not an exhaustive study of all possible error sources, it gives a good picture of the main challenges to make a precise and accurate α measurement with a BEC contrast interferometer. With the scattering length being tuned to zero, the main error contributions come from technical sources, especially the laser beams creating the standing wave. If the laser intensity could be stabilized to 0.2 %, the statistical error could be expected to be around 23 ppb per shot. Since it takes about 10 s to create a BEC in our setup, a sensitivity of 5 ppb would be attained in \approx 20 shots or 3.5 minutes, and 1 ppb precision could be reached in 1.5 h. The main systematic error would come from the Gouy phase shift. Although a competitive

4. Bloch oscillations and interaction-induced dephasing

measurement of α would be a difficult and time-consuming task requiring a purpose-built apparatus, we intend to make a proof-of-principle measurement with the current setup. A laser system capable of doing the required Bragg and Kapitza-Dirac pulses has already been built [Mar07a].

It is very clear that for atoms without a tunable scattering length, errors due to the meanfield energy dwarf all other errors. With a scattering length that is precisely tunable close to zero, the mean-field shifts would not be the primary issue any more. This is one of the main messages of this thesis.

5

Coherent dephasing of Bloch oscillations

In the previous chapter, we saw that atom-atom interactions lead to dephasing of Bloch oscillations. We will now demonstrate that, in the limit of a strong applied force, this dephasing causes the condensate wave function to develop an interference pattern within the first Brillouin zone. In contrast to other matter wave interference experiments, where particleparticle interactions are expected to lead to uncontrollable phase shifts and form a source of decoherence, the dynamics here are driven by the interactions. This illustrates an important point: this dephasing is a coherent process that can be controlled and even reversed using an external potential or a Feshbach resonance.

5.1 Theory

In this section, a simple analytic model for the dephasing of Bloch oscillations will be developed. It is a simplified version of a model described in [Wit05].

5.1.1 Dephasing in the limit of a strong force

In section 3.1.4, it was shown that the evolution of a BEC in a 1D optical lattice subject to an external force can be described by a discrete nonlinear Schrödinger equation of the form

$$i\hbar\frac{\partial c_j}{\partial t} = J\left(c_{j-1} + c_{j+1}\right) + \left(\mu_j^{loc}(c_j) + V_j^{trap} + Fdj\right)c_j,\tag{5.1}$$

where $c_j = \sqrt{N_j}e^{i\phi_j}$ is the complex amplitude of the wave function at lattice site *j*. In the case where the force *F* is large, this model can be further simplified. When the difference in potential energy between neighboring lattice sites, *Fd*, is much larger than the tunneling matrix element *J*, tunneling between lattice sites is strongly suppressed. The hopping term coupling adjacent lattice sites can then be neglected and the number of atoms at a site is fixed to its initial value. The time dependence of the system is now only determined by the phase

Coherent dephasing of Bloch oscillations

evolution of the local wave functions, and equation (5.1) can be rewritten as

$$\hbar \frac{d\phi_j}{dt} = \mu_j^{loc}(c_j) + V_j^{trap} + F dj.$$
(5.2)

As seen in equation (3.34), the local chemical potential μ_j^{loc} can be calculated from the density distribution as

$$\mu_j^{loc} = U_1 |c_j| = \sqrt{\frac{m\omega_\perp g}{\sqrt{2\pi}\pi\sigma_{\text{lat}}}} |c_j|.$$
(5.3)

When the lattice is adiabatically loaded, μ_j^{loc} can also be calculated in a much simpler way. In the Thomas-Fermi approximation of the BEC ground state (see section 3.1.5), the local chemical potential mirrors the external trapping potential,

$$\mu_j^{loc} = \mu - V_j^{trap}.\tag{5.4}$$

This has one important consequence; contrary to what one could believe from equation (5.3), the local chemical potential is *not* dependent on the scattering length used at loading, as long as the lattice loading is sufficiently adiabatic and the condensate stays in the ground state. On the other hand, after an external force is applied and the density distribution is frozen, the local chemical potential can be changed according to equation (5.3) by, for example, changing the scattering length.



Figure 5.1: If the lattice is adiabatically loaded and the condensate stays in the ground state, the local chemical potential mirrors the external trapping potential.

Knowing μ_j^{loc} , we can now calculate the BEC wave function for our system. The lattice is adiabatically loaded, the chemical potential mirrors the trap potential at loading and we therefore have $\mu_j^{loc} = \mu - \frac{1}{2}m\omega_z^2 d^2 j^2$. Bloch oscillations are induced by applying a force and shutting the dipole trap off, and we have $V_j^{trap} = 0$. The phase at lattice site j will then start to evolve as

$$\hbar \frac{d\phi_j}{dt} = -\alpha_{\rm int} j^2 + F dj, \qquad (5.5)$$

where $\alpha_{int} = \frac{1}{2}m\omega_z^2 d^2$ and a global phase due to μ has been neglected. Following the discussion in section 4.1, we can then write the wave function in momentum space as

$$\Phi(p,t) = \sum_{j} c_{j} w_{j}(p) = w_{0}(p) \sum_{j} c_{j} e^{-ipkd/\hbar}.$$
(5.6)

Using $c_j(t) = c_j(0)e^{\phi_j(t)}$ and converting to quasimomentum space by dividing by the Wannier function, we have

Ō

$$\Phi(q,t) = \sum_{j} c_j(0) e^{i(-\alpha_{\rm int}j^2 + Fdj)t/\hbar} e^{-ijqd/\hbar}$$
$$= \sum_{j} c_j(0) e^{-i\alpha_{\rm int}j^2t/\hbar} e^{-i(q-Ft)jd/\hbar}.$$
(5.7)

As previously discussed, the term linear in j describes a translation in quasimomentum space, causing Bloch oscillations. Now there is also a new term quadratic in j, coming from the spatially quadratic interaction potential, which leads to a dephasing between lattice sites. How does this affect the quasimomentum distribution? As visualized in figure 5.2a, the quadratic dephasing initially leads to a broadening of the Bloch oscillating wave function. Figure 5.2b shows how the wave function behaves for longer evolution times (here, the linear term leading to Bloch oscillations has been omitted, to better visualize the effect of the quadratic dephasing). We see that the wave function does not only broaden, but the interference of the wave functions at the different lattice sites gives rise to a complex pattern.



Figure 5.2: (a) The evolution of the quasimomentum distribution of a Bloch oscillating interacting BEC according to the model of equation (5.7) **(b)** Effect of the interaction term on the quasimomentum distribution for longer time scales.

5.1.2 Shift in oscillation period due to the lattice position

In the derivation of equation (5.7), it was assumed that the local chemical potential is of the form $\mu_j^{loc} = \alpha j^2$. However, if the center of the external harmonic confinement is not directly over one lattice site, this has to be slightly modified to $\mu_j^{loc} = \alpha (j + \delta)^2$, where $\delta \in [0, 1]$ describes an offset of the trap center with respect to one of the lattice minima. Equation (5.5) for the phase evolution then has to be modified to

$$\hbar \frac{d\phi}{dt} = F dj + \alpha (j + \delta)^2$$

= $(Fd + 2\alpha\delta)j + \alpha j^2 + const.$ (5.8)

5. Coherent dephasing of Bloch oscillations

The offset results in an additional term linear in j and can be seen as adding a modification to the applied force F. In the experiment, we have $Fd \gg 2\alpha$ and the modification is small, but not negligible. In practice, it is very challenging to stabilize the external trap position with respect to the location of the lattice wells. This means that the quasimomentum distribution of two subsequent experimental realizations will look equal, except that the center of the distribution will be shifted to due to the effective force being slightly different.

5.1.3 From the 1D model to a 2D absorption picture

Absorption pictures taken in our experiment after a long expansion time show the momentum distribution of the wave function integrated over one direction,

$$F(p_x, p_z) = \int |\Psi(\boldsymbol{p})|^2 dp_y.$$
(5.9)

To compare our simple model with the absorption pictures, we need to calculate this quantity. The total 3D wave function in momentum space can be written as a sum of a Wannier function in the vertical (*z*-) direction and a Thomas-Fermi wave function in the radial (ρ -) direction,

$$\Psi(\boldsymbol{p},t) = \sum_{j} c_j(t) e^{-ijp_z d/\hbar} \Phi_W(p_z) \psi_{TF}(p_\rho, N_j).$$
(5.10)

The Thomas-Fermi wave functions ψ_{TF} are dependent on the atom number and equation (5.10) cannot be directly simplified. It is however straightforward to evaluate equations (5.9) and (5.10) numerically.

5.2 Experimental realization

5.2.1 Experimental parameters

The setup for studying the coherent dephasing is essentially identical to the setup described in chapter 4. The only substantial difference is that we work with a dipole trap that is about 4 times tighter in the vertical direction, which makes the dephasing faster.

The experiments are performed with a BEC of up to $1.5 \cdot 10^5$ atoms created in a dipole trap with trap frequencies $2\pi \times (39, 5, 39)$ Hz. The scattering length is set to 210 a_0 . The trap is tightened to $2\pi \times (41, 13, 39)$ Hz and at the same time the optical lattice oriented along the vertical direction is raised to 8 E_R within 1 s. The slow ramp makes sure that the BEC is loaded adiabatically in the lattice, even with respect to transversal breathing modes. After loading, the BEC is cigar-shaped with Thomas-Fermi radii (9,28,8) µm. This means that approximately 35 lattice sites are occupied, with up to 7500 atoms in the central site.

As before, Bloch oscillations are induced by turning off the magnetic levitation while simultaneously ramping down the power in the horizontal dimple beam in 0.3 ms. The coils producing the levitation gradient also add a small homogenous magnetic field. Therefore, shutting off these coils slightly changes the magnetic field at the position of the atomic sample, and the scattering length is changed to 190 a_0 . After a period of evolution in the lattice, the vertical dimple beam and the lattice are adiabatically ramped down in 1 ms.



Figure 5.3: Interaction-induced matter wave interference. (a) Absorption images showing the quasimomentum distribution as a function of time. (b) The evolution of the quasimomentum wave function evolution according to the simple model in equation (5.7). The pictures are simulated absorption images calculated as described in section 5.1.3. α_{int} has been rescaled by a factor 0.9 to account for the reduction in density due to the transversal dynamics, see text.

5.2.2 Structure of a dephased cloud in quasimomentum space

According to the simple model discussed in section 5.1.1, the dephasing of of the BEC due to interactions should give rise to interference patterns in quasimomentum space. We test this by taking absorption pictures of the dephased condensate after different times of evolution in the lattice. The pictures are taken after 80 ms of levitated expansion, where the vertical trapping beam was slowly ramped down over 50 ms to reduce the expansion in the horizon-tal direction. Figure 5.3a shows the result of this experiment. During the first couple of Bloch cycles, the wave function broadens and spreads out in the first Brillouin zone. After about 14 Bloch cycles, interference fringes start to develop at the edge of the wave function. The fringes later grow inwards, and after about 26 cycles cover the whole Brillouin zone. The evolution of the pattern can be followed over more than 100 Bloch cycles, corresponding to over 60 ms.

We find that the distance between the interference maxima is reproducible, but the overall position of the pattern will be slightly shifted from one experimental run to the next. This can be attributed to the subtle effect described in the previous chapter; the exact location of the lattice minima with respect to the dipole trap minimum will cause a slight change of the Bloch frequency.

The images obtained in the experiment can be compared with the simple analytical model from section 5.1.1. For our experimental parameters, $Fd \approx 1.3E_R \approx h \times 1740$ Hz, $J \approx 0.038E_R \approx h \times 40$ Hz and $\alpha_{int} \approx 2.1 \cdot 10^{-3}E_R \approx h \cdot 2.8$ Hz. The evolution of the wave function is very well reproduced when we reduce α_{int} by a factor 0.9. This scale factor accounts primarily for the fact that our model assumes that there is no transversal dynamics. In the experiment, this is not the case. Turning off the horizontal trapping beam when the Bloch

5. Coherent dephasing of Bloch oscillations

oscillations start excites a radial breathing mode (there is no breathing along the vertical axis since tunneling between lattice wells is prohibited). The breathing reduces the density and modulates it in time. Reducing α_{int} is a simple way of taking this into account. Although not a complete description of the experiment, our simple model captures the essential physics: The mean-field interactions lead to a phase shift that is quadratic with respect to the lattice site number. As a result of this phase shift, the interference of the different lattice site wave functions leads to a momentum broadening and later an interference pattern. The emergence of an interference pattern and its agreement with the theoretical model demonstrates that the interaction-induced dephasing is a coherent process.



Figure 5.4: Left: Contrast of the interference fringes emerging during free expansion after the trap is switched off. The BEC has been dephased over 40 Bloch cycles. The contrast is defined as $(I_{max} - I_{min}) / (I_{max} + I_{min})$, where $I_{max}(I_{min})$ is the average of the maxima (minima) of the central peak structure. **Right:** Absorption pictures and integrated column density for two different expansion times.

To observe the interference structure with high contrast, two key features of our imaging procedure are important. First, the scattering length has to be switched to near zero during expansion. This minimizes interactions during expansion, which otherwise lead to mean-field induced broadening of the cloud. Second, the time of free expansion must be long, such that the size of the imaged cloud is much larger than the size of the condensate in the trap, and the absorption image really shows the momentum distribution. This is illustrated in figure 5.4. However, if the expansion time is too long, the atom cloud becomes very dilute and the signal-to-noise ratio goes down. We found 80 ms expansion time to be a good compromise, allowing both high contrast fringes and low-noise pictures. Throughout this chapter, this is the expansion time used unless otherwise indicated.

5.2.3 Cancelation of dephasing through an external potential

The dephasing of the BEC in our system is due to the interaction potential. Equation (5.2) suggests that it should be possible to counteract the interaction potential using a suitable external potential. We can easily test this by switching the dipole trap depth to a new value during the Bloch oscillations, instead of shutting the trap off as in the previous case. The result of this experiment can be seen in figure 5.5. Figure 5.5a and 5.5b show absorption pictures taken after 40 Bloch oscillations. Without a compensating potential, the BEC has

dephased and the interference structure can be seen. However, with the right strength of the compensating potential, Bloch oscillations are still clearly visible. This is further quantified in figure 5.5c, which plots the rms-width of the quasimomentum distribution. The width is minimized if the laser power of the trap is close to the one used at loading. This means that, in contrast to most previously performed experiments, the best strategy for seeing Bloch oscillations in a system with interactions is to not turn off the trap and instead only apply a constant force.



Figure 5.5: Absorption images showing the quasimomentum distribution after 40 Bloch oscillations for phase π (a) and phase 0 (b) and the quasimomentum rms-width (c) as a function of the external confinement strength, normalized to the strength after loading. (d) Evolution of the momentum distribution for the case when no compensating potential is present, showing fast broadening. (e) When the effect of the mean-field potential is canceled by an external potential, Bloch oscillations are visible for a much longer time.

This effect is straightforward to understand. When the lattice is loaded, the atoms distribute themselves such that the local chemical potential is the same at every lattice site, as illustrated in figure 5.6. If the trap is turned off when the Bloch oscillations are started, both a linear (from the force) and a quadratic (from the atom-atom interactions) potential is imparted on the atoms. By keeping the trap on, only a linear potential is applied which leads to clean Bloch oscillations.

For how long can one observe Bloch oscillations when the interaction potential is optimally compensated? Figure 5.5e shows how the quasimomentum distribution evolves with time when interactions are optimally compensated. We see that the Bloch oscillations are visible for about 150 Bloch cycles. This is an order of magnitude longer than for the case without a compensating potential, shown in figure 5.5d. Here, after about 15 Bloch cycles, the wave function is already so broad that it is hard to follow the Bloch oscillations. It is also interesting to study the way the wave function dephases when the compensating potential is added. Instead of broadening, the width of the central peak initially stays the same. Over time, the peak is slowly depopulated, while the population in a broad background distribution increases. After about 100 Bloch cycles, the central peak starts to develop side lobes or splits in two, with the exact shape varying from one experimental run to the next.

5. Coherent dephasing of Bloch oscillations



Figure 5.6: Illustration of the local chemical potential during Bloch oscillations. Left: When the lattice is loaded, the BEC arranges itself such that the local chemical potential is constant. Middle: The trap is turned off when Bloch oscillations are started. The local chemical potential gets a parabolic shape, leading to dephasing. Right: The dipole trap is not turned off, and the local chemical potential varies linearly, leading to Bloch oscillations without dephasing.

Other sources of dephasing

According to our simple model, there should be no dephasing when the effect of interactions is compensated with the external potential. What effects not included in the model limit the number of Bloch oscillations one can observe? I will discuss several possibilities: Radial breathing, dynamical instabilities and beyond-mean-field effects leading to decoherence of the phase on a single lattice site.

Radial breathing

Our simple model assumes that there is no motion in the radial direction. However, radial breathing modes are excited when the Bloch oscillations are started, especially in the case where the external harmonic trapping potential is turned off, leading to a change of the external confinement. The radial breathing, modulating the density and therefore the interaction energy, could couple the different degrees of freedom and lead to additional dephasing. When the interaction energy is compensated by the external potential as in figure 5.5e, the external confinement is not changed when the Bloch oscillations are started, and only a small radial excitation due to the magnetic field switching is visible. It is therefore unlikely that this is the source for the residual dephasing observed.

Dynamical instabilities

Not all solutions to the Gross-Pitaevskii equation are stable. There are cases where the eigenfrequencies of the Bogoliubov modes of the system exhibit complex values. Small perturbations of the wave function grow exponentially over time, and these excitations ultimately lead to heating and destruction of the condensate. This kind of instability is called a dynamical instability.

When no external force is present, Bloch states that are close to the edge of the Brillouin zone are dynamically unstable [Wu01, Wu03, Fal04, DS05]. Does this picture still hold when an external force is applied? We can examine this effect by inducing Bloch oscillations as before, but vary the external force by applying different levitating magnetic field gradients.

Figure 5.7 shows an example of Bloch oscillations with a small force. With an acceleration of about 0.07 g, coherence is destroyed as soon as the edge of the Brillouin zone is reached and the distribution in quasimomentum space becomes evenly distributed. For a larger acceleration of 0.14 g, the condensate partly "survives" the first crossing of the Brillouin zone edge, but after the second crossing coherence is completely lost. This agrees with the findings in [Cri04], where the first crossing of the Brillouin zone edge was investigated for a few different lattice accelerations.



Figure 5.7: Bloch oscillations of an interacting BEC with a small external force, corresponding to an acceleration of 0.07 g and 0.14 g, respectively. The scattering length is about 190 a_0 .

In [Cri04], the interpretation of this result was that with a larger acceleration, the condensate spends less time in the unstable region at the edge of the Brillouin zone, and therefore excitations do not have as much time to grow. However, stability analysis of a simple 1D model [Zhe04] indicates that this picture may be too simple. Instead, if the applied force is large enough that the potential energy difference between two adjacent lattice sites is large compared to the interactions, characterized by the local chemical potential, the condensate is stable in all regions of the Brillouin zone. The stability criterion is

$$Fd > 3.03\mu^{loc}$$
. (5.11)

If the force is smaller than the critical force $F_c = 3.03 \mu^{loc}/d$, the tunneling strength becomes important. The stability criterion is then

$$Fd/\mu^{loc} > 4.186\sqrt{J/\mu^{loc}}.$$
 (5.12)

The region of stability is plotted in figure 5.8. In the experiments presented in this chapter, with 1*g* acceleration and a lattice depth of 8 E_R , we have $Fd/\mu^{loc} \approx 2.38$ and $\sqrt{J/\mu^{loc}} \approx 0.24$. This point is marked in the figure, and it is within the stable region.

Figure 5.9 shows the condensate quasimomentum wave function after forces of different magnitude have been applied for a fixed amount of time, 30 ms. For this evolution time, the

5. Coherent dephasing of Bloch oscillations



Figure 5.8: Regions of stability for a Bloch oscillating BEC. The vertical line marks the critical force $F_c = 3.03 \mu^{loc}/d$ and the diagonal line is $Fd = 4.186 \sqrt{J/\mu^{loc}}$. The cross marks the position corresponding to the parameters used in the experiments described in this chapter when the acceleration is 1 g.

condensate has had time to undergo many Bloch oscillations and for both the smaller and the larger accelerations, approximately the same time has been spent crossing the edge of the Brillouin zone. However, for accelerations $\leq 0.4 g$, the atoms are homogeneously distributed over the Brillouin zone. For larger accelerations the interference pattern appears, a sign of coherent evolution. This is consistent with the model put forward in [Zhe04].

Beyond-mean-field effects, on-site phase diffusion

Until now, we have assumed that we are in a regime where mean-field theory is applicable, and the system can be described in a Gross-Pitaevskii framework. Going beyond the mean-field treatment, several factors can lead to dephasing. For example, the number of atoms at each lattice site is not a fixed number, but the many-body ground state exhibits number fluctuations. For low lattice depths, where $\mu^{loc}/J \ll 1$, the wave function at an individual lattice site can be constructed as a superposition of different number states as $|\Psi\rangle = \sum f(n) |n\rangle$ with the distribution f(n) being Poissonian with a width $\sigma(N) = \sqrt{N}$, where N is the atom number [Jav99].

When the Bloch oscillations are started by a applying a potential gradient, tunneling is suppressed and the lattice sites become isolated from each other. The many-body wave function can be written as a product state of the wave functions at the individual lattice sites, which evolve independent of each other. The single-site contribution to the wave function will evolve as

$$|\Psi(t)\rangle = \sum_{n} f(n)e^{-iE(n)t/\hbar} |n\rangle, \qquad (5.13)$$

where E(n) is the energy corresponding to the number state $|n\rangle$. This means that the initially

5.2. Experimental realization



Figure 5.9: Absorption pictures showing the quasimomentum distribution after an external force was applied for 30 ms. In the first picture, the force is so small that the edge of the Brillouin zone has not yet been reached, and the condensate is still intact. If the force is increased, the condensate has passed the edge of the Brillouin zone and is destroyed by a dynamical instability. However, above approximately 0.4 g acceleration, corresponding to $Fd/\mu^{loc} \approx 4\sqrt{J/\mu^{loc}}$, a clear interference pattern is visible, indicating the absence of dynamical instabilities. In this region the dephasing is coherent.

well-defined phase will diffuse with time. In the limit of large atom numbers E(n) can be linearized, and it can be shown that the dephasing then happens with a decoherence time [Ima97, Li07].

$$\tau_{coh} = \frac{\hbar N}{\mu^{loc}\sigma(N)} = \frac{\hbar\sqrt{N}}{\mu^{loc}}.$$
(5.14)

In our system each lattice site contains a 2D condensate, and the local chemical potential at a lattice site then is $\mu^{loc}(N) = U_1 \sqrt{N}$, with U_1 defined in equation (3.34). This leads to the surprising consequence that the decoherence time becomes independent of atom number and simply reads $\tau_{coh} = h/U_1$.

However, with a deeper lattice the relation $\sigma(N) = \sqrt{N}$ no longer holds. The combination of interactions and decreased tunneling make number fluctuations energetically unfavourable. The uncertainty in atom number becomes sub-Poissonian [Jav99],

$$\sigma(N) = \sqrt{N} \left(\frac{1}{1 + \frac{\mu^{loc}}{2J}}\right)^{1/4}.$$
(5.15)

This is called *number squeezing*. Using the same reasoning as above, the coherence time now becomes

$$\tau_{coh} = \frac{h\sqrt{N}}{\mu^{loc}} \left(1 + \frac{\mu^{loc}}{2J}\right)^{1/4} = \frac{h}{U_1} \left(1 + \frac{U_1\sqrt{N}}{2J}\right)^{1/4}.$$
(5.16)

Note that even with the correction due to number squeezing, the coherence time depends very weakly on the atom number. In the limit of large N, it scales as $N^{1/8}$. However, in our system, the correction due to number squeezing is still important. The decoherence time is 20 ms for our experimental parameters, a substantial difference to the 10 ms decoherence time one calculates when not taking number squeezing into account.

The decoherence time we calculate is on the same order of magnitude as the time it takes for the wave function in our system to dephase, as seen in figure 5.5. We believe that the on-site phase diffusion and other beyond-mean-field effects, such as the interaction with a remnant thermal cloud, cause the residual dephasing we observe in our experiment. 5. Coherent dephasing of Bloch oscillations

5.2.4 Rephasing of a dephased condensate

Another interesting demonstration of the coherence in our system is to rephase a dephased condensate, in the spirit of spin-echo techniques known from nuclear magnetic resonance experiments. The idea is to turn off interactions and apply the external potential, thereby letting the lattice sites experience a potential with opposite sign, causing the phase evolution to "run backwards" until the the condensate is recovered. However, since a sudden switch of the scattering length excites very strong breathing in the horizontal direction, we employ a slightly modified protocol where the switching is more gentle and reduces breathing.

Using the same procedure to load the lattice and start the Bloch oscillations as in the previous experiments, we let the BEC Bloch oscillate with the dipole trap ramped down in 4 ms and the scattering length set to 214 a_0 for 35 Bloch cycles. The scattering length is then ramped to 10 a_0 within 10 ms. At the same time the dipole trap is ramped up in 4 ms to approximately the same depth as was used during loading.

As can be seen in figure 5.10, the wave function then refocuses and the initial BEC peak is recovered 24 Bloch cycles after the scattering length ramp. The shift in the Bloch frequency due to the trap positioning causes the quasimomentum distribution to shift around from one experimental run to the next. We recorded about 10 distributions for each evolution time and selected those which where symmetrical, corresponding to Bloch cycle phase $\phi = 0$ or $\phi = \pi$, where we define the phase of the Bloch oscillation such that $\phi = 0$ corresponds to the BEC wave function being centered in the middle of the Brillouin zone. As can be seen in the absorption pictures in the figure, there is still considerable horizontal breathing, but the rephasing nevertheless works very well. This further confirms that the initial broadening and dephasing has been coherent.

5.2.5 Decay and revival of Bloch oscillations

Until now, the dephasing of the Bloch oscillating condensate has been induced through interactions. We also have the possibility to switch off atom-atom interactions and "artificially" induce dephasing through an external potential, the dipole trap. The decoherence mechanisms listed in section 5.2.3 no longer apply, and our simple analytical model is valid over a long time. This allows not only to observe the dephasing, but also to see a revival of the Bloch oscillations.

With interactions turned off and a dipole trap with trap frequency ω_z along the lattice direction, the phase evolution at lattice site *j* is

$$\hbar\phi_j = Fdjt + \frac{1}{2}m\omega_z^2 d^2 j^2 t.$$
 (5.17)

If we wait long enough, the dephasing term will be a multiple of 2π for all j - the lattice sites have come back in phase, the original wave function is recovered and Bloch oscillations can again be seen [Pon06]. This happens when $\frac{1}{2}m\omega_z^2 d^2t/\hbar = 2\pi$. Less obvious is that the original wave function also will be recovered at the earlier time T_R when $\frac{1}{2}m\omega_z^2 d^2T_R/\hbar = \pi$. This can be shown with a simple calculation. The wave function after a time T_R is

$$\Psi(q, T_R) = \sum_j c_j(0) e^{-i\pi j^2} e^{-iqjd/\hbar}$$
(5.18)



Figure 5.10: A fully dephased BEC is rephased back into a narrow momentum distribution by switching interactions close to zero and applying an external harmonic potential, rewinding the phase at the different lattice sites. The black solid lines show quasimomentum distributions corresponding to Bloch cycle phase 0, separated in time by two Bloch oscillations or about 1.15 ms. The red solid lines correspond to quasimomentum distributions half a Bloch oscillation later (cycle phase π). The images are absorption images corresponding to the adjacent quasimomentum distributions. Note the excitation along the radial direction in the images.

5. Coherent dephasing of Bloch oscillations

(for simplicity, the phase term leading to Bloch oscillations has been omitted). Using the relation $e^{-i\pi j^2} = (-1)^{j^2} = (-1)^j = e^{-i\pi j}$, this can be rewritten as

$$\Psi(q, T_R) = \sum_j c_j(0) e^{-i\pi j} e^{-iqjd/\hbar}$$
$$= \sum_j c_j(0) e^{-i(q+\hbar k)jd/\hbar}$$
$$= \Psi(q + \hbar k, 0).$$
(5.19)

We see that the original wave function is recovered, only with the quasimomentum shifted by $\hbar k$. This means that we would expect a revival of the dephased Bloch oscillations after the time

$$T_R = \frac{h}{m\omega_z^2 d^2}.$$
(5.20)

Expressed in number of Bloch oscillations, the value is

$$N_R = \frac{F}{m\omega_z^2 d} = \frac{g}{\omega_z^2 d}.$$
(5.21)



Figure 5.11: Decay and subsequent revival of Bloch oscillations. The rms-width of the momentum distribution is plotted as a function of the number of Bloch oscillations for Bloch cycle phase 0 (white circles) and π (black circles). Absorption images corresponding to N = 1, 70, 140, 210 and 280 Bloch oscillations are also shown.

We perform an experiment with a non-interacting Bloch oscillating BEC where a dipole trap with trap frequency $\omega_z = 40(1)$ Hz is applied. The result is shown in figure 5.11. The BEC is quickly dephased, but after about 280 Bloch cycles, Bloch oscillations are again visible. This is consistent with the expected value of $N_R = 292(15)$ Bloch cycles. Note that after half the revival time, two clouds can be seen Bloch oscillating. It can be shown that after a time T_R/k , k copies of the original wave function will appear. This is called the fractional Talbot effect [Ber96, Kap00].

The uncertainty in dipole trap position relative to the lattice sites, which slightly affects the Bloch period (see section 5.1.2), becomes increasingly important for longer evolution

times. Using equation (5.8), one finds that the positioning of the dipole trap adds an extra phase factor which at the time of the revival is $2\pi\delta$. Since we cannot control the trap positioning, the phase of the Bloch oscillation shifts from one experimental realization to the next, and the data in figure 5.11 has been obtained by acquiring many images and selecting those corresponding to cycle phase 0 or π .

The phase of the Bloch oscillation at the time of the first revival gives information on the trap position relative to the lattice with a precision only limited by how precisely the phase can be measured. This means that one can measure the trap positioning to a fraction of a wavelength! With active stabilization of the dipole trap beam pointing and interferometric stabilization of the lattice phase, one could investigate this effect in more detail. This would be an interesting endeavour.

6 Outlook

In the framework of this thesis, an apparatus allowing experiments with a tunable quantum gas in an optical lattice was developed. We have demonstrated a precise control over atom-atom interactions, and this work opens up a whole realm of opportunities for future experiments. I will here outline a few interesting possibilities for future work, some of which have already been implemented at the time of writing.

Quantum gases of deeply bound molecules

Cs has a very rich molecular structure and a quantum gas of Cs₂ molecules can be created out of a Cs BEC by sweeping the magnetic field over a Feshbach resonance [Her03]. These weakly bound "Feshbach molecules" are not collisionally stable and one attractive goal has been to find an efficient way of populating deeply bound states, especially the rovibronic ground state, which should be stable against two-body loss. Coherent transfer from one molecular state to another can be accomplished using a two-photon transition and the stimulated adiabatic rapid passage (STIRAP) technique. A transfer from the most weakly bound vibrational level to the next bound level, with a binding energy of 640 MHz, was demonstrated last year in the Innsbruck Rb experiment [Win07b].

A transfer to a more deeply bound molecular state poses two major challenges. First, the deeply bound target state has a binding energy on the order of 140 THz. The STIRAP technique requires the two transfer lasers to be phase coherent, despite being detuned from each other by many THz. Second, the weakly bound starting state has a large spatial extent and the wave function overlap with the spatially small target state is low. We managed to overcome these problems by performing spectroscopy to identify levels with adequate transition rates [Dan08a] and by referencing the transfer lasers to a frequency comb. As illustrated in figure 6.1, we were then able to perform a STIRAP transfer of Feshbach molecules into the $|v = 73, J = 2\rangle$ state, a state bound by 32 THz, with about 80 % efficiency and no discernible heating [Dan08b]. With the addition of two more transfer lasers, a further STIRAP step makes it possible to reach the rovibronic ground state. We have identified a suitable intermediate state for the transfer [Mar08], and preliminary experiments have reached a transfer efficiency for the second step of about 75 %.

The availability of an optical lattice in the apparatus could help increase the efficiency of



Figure 6.1: Illustration of STIRAP transfer of Cs₂ molecules from a weakly bound state to a deeply bound state, adapted from [Dan08b]. (a) Molecular level scheme for Cs₂. Molecules are transfered from the weakly bound vibrational level $|v = 155\rangle$ of the singlet $X^1\Sigma_g^+$ potential to the deeply bound $|v = 73\rangle$ level using a two-photon STIRAP process, with level $|v' \approx 225\rangle$ of the electronically excited 0_u^+ potential as an intermediate transfer state. (b) Number of $|v = 155\rangle$ molecules detected as a function of STIRAP time τ . After about 15 µs, all Feshbach molecules have been transfered to $|v = 73, J = 2\rangle$, a state which cannot be directly detected. After 30 µs, 65 % of the molecules have been transferred back into $|v = 155\rangle$, indicating a single pass transfer efficiency of 80 %. (c) Schematic timing of the STIRAP transfer process. (d) Round-trip transfer efficiency as a function of detuning Δ_2 of transfer laser L₂, with laser L₁ on resonance.

the initial Feshbach molecule production. The creation of a Mott insulator state [Jak98, Gre02] with two atoms per lattice site allows production of Feshbach molecules with high efficiency and with strongly suppressed collisional losses [Tha06, Vol06]. A high number of Feshbach molecules would be an excellent starting point for the creation of a BEC of ground state molecules.

One-dimensional systems, Tonks-Girardeau gas

A two-dimensional optical lattice can be raised to create an array of tight tube-shaped traps. If the lattice is deep enough that the vibrational level spacing in the tight direction is much larger than the temperature and the chemical potential of the atomic sample, radial excitations are "frozen out" and the dynamics is one-dimensional. A key parameter in such a quasi-1D system is γ , the ratio of interaction energy to kinetic energy [Pet00]. In the weakly interacting regime where $\gamma \ll 1$, a quantum degenerate system is a Bose-Einstein condensate. The regime where $\gamma \gg 1$ is called the Tonks-Girardeau regime, where the system can be described as a gas of point-like impenetrable bosons and acquires fermionic properties [Gir60]. The ability to tune the scattering length makes it possible to reach deep into this regime, and at the time of writing, a Tonks-Girardeau gas with $\gamma \approx 50$ has been realized in our lab. We have investigated how the frequency of the lowest compressional mode of the gas varies with the scattering length, which also allowed us to reach beyond the Tonks-Girardeau regime and demonstrate a so-called super-Tonks gas [Ast05]. In the future, the ability to add a further lattice along the tubes makes it possible to investigate the commensurate-incommensurate transition [Büc03a].

Contrast interferometer at zero scattering length

The technique to strongly suppress atom-atom interactions demonstrated in this thesis has interesting applications in atom interferometry. The high brightness of a BEC combined with its narrow spatial and momentum spread would constitute an ideal source for a matter wave interferometer. Unfortunately, because of the high density in a BEC, interactions lead to phase diffusion and can cause systematic frequency shifts due to unwanted density gradients, limiting the performance of the interferometer. This limitation could be overcome using the precise control over the interaction strength demonstrated in this work.

As discussed in section 4.3, a BEC contrast interferometer can be used to measure the photon recoil frequency for Cs, from which the fine structure constant α can be determined. A previous implementation [Gup02] was severely limited in sensitivity and accuracy by interaction shifts. The laser setup needed to implement such a contrast interferometer has already been implemented [Mar07a] and we plan to do a proof-of-principle experiment demonstrating the suppression of interaction shifts in the near future.

Three-body loss in an optical lattice

In an gas with density n, the two-body collision rate is proportional to $\langle n^2 \rangle$. When a BEC is loaded into a 3D optical lattice, the high peak densities at the lattice sites lead to an effective increase in the two-body collision rate, as discussed in section 3.1.5. The number of threebody collisions, proportional to $\langle n^3 \rangle$, are even more increased. Combined with the intrinsic high three-body loss in Cs, this makes it possible to achieve very large three-body loss rates.

6. Outlook

Paradoxically, large three-body loss rates are expected to decrease the atom loss from the lattice, because tunneling to create a triply occupied lattice site is strongly suppressed, a form of quantum Zeno effect. The strong suppression of triply occupied lattice sites can be viewed as an effective large three-body interaction, which makes it possible to engineer new quantum phases [Dal08].

Bibliography

- [And95] M. H. Anderson, J. R. Ensher, M. R. Matthews, C. E. Wieman, and E. A. Cornell, Observation of Bose-Einstein Condensation in a Dilute Atomic Vapor, Science 269, 198 (1995).
- [And97] M. R. Andrews, C. G. Townsend, H.-J. Miesner, D. S. Durfee, D. M. Kurn, and W. Ketterle, Observation of Interference Between Two Bose Condensates, Science 275, 637 (1997).
- [And98] B. P. Anderson and M. A. Kasevich, *Macroscopic Quantum Interference from Atomic Tunnel Arrays*, Science **282**, 1686 (1998).
- [Arl98] J. Arlt, P. Bance, S. Hopkins, J. Martin, S. Webster, A. Wilson, K. Zetie, and C. J. Foot, Suppression of collisional loss from a magnetic trap, J. Phys. B 31, L321 (1998).
- [AS01] J. R. Abo-Shaeer, C. Raman, J. M. Vogels, and W. Ketterle, *Observation of Vortex Lattices in Bose-Einstein Condensates*, Science **292**, 476 (2001).
- [Ash76] N. W. Ashcroft and D. N. Mermin, *Solid State Physics*, Saunders College Publishing, 1976.
- [Ast05] G. E. Astrakharchik, J. Boronat, J. Casulleras, and S. Giorgini, *Beyond the Tonks-Girardeau Gas: Strongly Correlated Regime in Quasi-One-Dimensional Bose Gases*, Phys. Rev. Lett. **95**, 190407 (2005).
- [Aud03] G. Audi, A. H. Wapstra, and C. Thibault, *The 2003 atomic mass evaluation: (II). Tables, graphs and references,* Nucl. Phys. A **729**, 337 (2003).
- [BD96] M. Ben Dahan, E. Peik, J. Reichel, Y. Castin, and C. Salomon, *Bloch Oscillations of Atoms in an Optical Potential*, Phys. Rev. Lett. **76**, 4508 (1996).
- [Ber96] M. Berry and S. Klein, *Integer, fractional and fractal Talbot effect*, J. Mod. Opt. **43**, 2139 (1996).
- [Blo29] F. Bloch, Über die Quantenmechanik der Elektronen in Kristallgittern, Z. Phys. **52**, 555 (1929).
- [Blo08] I. Bloch, J. Dalibard, and W. Zwerger, Many-body physics with ultracold gases, Rev. Mod. Phys. 80, 885 (2008).
- [Bos24] S. Bose, *Plancks Gesetz und Lichtquantenhypothese*, Z. Phys. 26, 178 (1924).

- [Büc03a] H. P. Büchler, G. Blatter, and W. Zwerger, *Commensurate-Incommensurate Transition* of Cold Atoms in an Optical Lattice, Phys. Rev. Lett. **90**, 130401 (2003).
- [Büc03b] M. Büchner, R. Delhuille, A. Miffre, C. Robilliard, J. Vigué, and C. Champenois, *Diffraction phases in atom interferometers*, Phys. Rev. A **68**, 013607 (2003).
- [Cas01] Y. Castin, Bose-Einstein condensates in atomic gases: simple theoretical results, in: R. Kaiser, C. Westbrook, and F. David (Eds.), 'Coherent atomic matter waves', Lecture Notes of Les Houches Summer School, 1–136, EDP Sciences and Springer-Verlag, 2001, arXiv:cond-mat/0105058.
- [Chi00] C. Chin, V. Vuletic, A. J. Kerman, and S. Chu, *High Resolution Feshbach Spectroscopy* of *Cesium*, Phys. Rev. Lett. **85**, 2717 (2000).
- [Chi01] C. Chin, *Cooling*, *Collisions and Coherence of Cold Cesium Atoms in a Trap*, Ph.D. thesis, Stanford University (2001).
- [Chi04] C. Chin, V. Vuletić, A. J. Kerman, S. Chu, E. Tiesinga, P. J. Leo, and C. J. Williams, *Precision Feshbach spectroscopy of ultracold Cs*₂, Phys. Rev. A **70**, 032701 (2004).
- [Cla05] P. Cladé, S. Guellati-Khélifa, C. Schwob, F. Nez, L. Julien, and F. Biraben, Noise sensitivity of an atomic velocity sensor, Eur. Phys. Jour. D 33, 173 (2005).
- [Cla06a] P. Cladé, E. de Mirandes, M. Cadoret, S. Guellati-Khélifa, C. Schwob, F. Nez, L. Julien, and F. Biraben, *Determination of the Fine Structure Constant Based on Bloch Oscillations of Ultracold Atoms in a Vertical Optical Lattice*, Phys. Rev. Lett. 96, 033001 (2006).
- [Cla06b] P. Cladé, E. de Mirandes, M. Cadoret, S. Guellati-Khélifa, C. Schwob, F. Nez, L. Julien, and F. Biraben, *Precise measurement of h/m_{Rb} using Bloch oscillations in a vertical optical lattice: Determination of the fine-structure constant*, Phys. Rev. A **74**, 052109 (2006).
- [Cri02] M. Cristiani, O. Morsch, J. H. Müller, D. Ciampini, and E. Arimondo, Experimental properties of Bose-Einstein condensates in one-dimensional optical lattices: Bloch oscillations, Landau-Zener tunneling, and mean-field effects, Phys. Rev. A 65, 063612 (2002).
- [Cri04] M. Cristiani, O. Morsch, N. Malossi, M. Jona-Lasinio, M. Anderlini, E. Courtade, and E. Arimondo, *Instabilities of a Bose-Einstein condensate in a periodic potential: an experimental investigation*, Opt. Express **12**, 4 (2004).
- [Dal99a] F. Dalfovo, S. Giorgini, L. P. Pitaevskii, and S. Stringari, *Theory of Bose-Einstein condensation in trapped gases*, Rev. Mod. Phys. 71, 463 (1999).
- [Dal99b] J. Dalibard, Collisional dynamics of ultra-cold atomic gases, in: M. Inguscio, S. Stringari, and C. E. Wieman (Eds.), Proceedings of the International School of Physics - Enrico Fermi, 321, IOS Press, 1999.
- [Dal08] A. J. Daley, J. M. Taylor, S. Diehl, M. Baranov, and P. Zoller, *Atomic three-body loss as a dynamical three-body interaction*, arXiv:0810.5153 (2008).

- [Dan08a] J. G. Danzl, E. Haller, M. Gustavsson, M. J. Mark, N. Bouloufa, O. Dulieu, H. Ritsch, R. Hart, and H.-C. Nägerl, *Precision molecular spectroscopy for ground* state transfer of molecular quantum gases, arXiv:0811.2374 (2008).
- [Dan08b] J. G. Danzl, E. Haller, M. Gustavsson, M. J. Mark, R. Hart, N. Bouloufa, O. Dulieu, H. Ritsch, and H.-C. Nägerl, *Quantum Gas of Deeply Bound Ground State Molecules*, Science **321**, 1062 (2008), published online 10 July 2008; 10.1126/science.1159909.
- [DD04] G. Di Domenico, N. Castagna, G. Mileti, P. Thomann, A. V. Taichenachev, and V. I. Yudin, Laser collimation of a continuous beam of cold atoms using Zeeman-shift degenerate-Raman-sideband cooling, Phys. Rev. A 69, 063403 (2004).
- [Den00] J. Denschlag, J. E. Simsarian, D. L. Feder, C. W. Clark, L. A. Collins, J. Cubizolles, L. Deng, E. W. Hagley, K. Helmerson, W. P. Reinhardt, S. L. Rolston, B. I. Schneider, and W. D. Phillips, *Generating Solitons by Phase Engineering of a Bose-Einstein Condensate*, Science 287, 97 (2000).
- [Den02] J. H. Denschlag, J. E. Simsarian, H. Haeffner, C. McKenzie, A. Browaeys, D. Cho, K. Helmerson, S. L. Rolston, and W. D. Phillips, A Bose-Einstein condensate in an optical lattice, J. Phys. B 35, 3095 (2002).
- [DeP00] M. T. DePue, S. L. Winoto, D. J. Han, and D. S. Weiss, *Transient compression of a MOT and high intensity fluorescent imaging of optically thick clouds of atoms*, Opt. Comm. 180, 73 (2000).
- [Don01] E. A. Donley, N. R. Clausen, S. L. Cornish, J. L. Roberts, E. A. Cornell, and C. E. Wieman, *Dynamics of Collapsing and Exploding Bose-Einstein Condensates*, Nature 412, 295 (2001).
- [DS05] L. De Sarlo, L. Fallani, J. E. Lye, M. Modugno, R. Saers, C. Fort, and M. Inguscio, Unstable regimes for a Bose-Einstein condensate in an optical lattice, Phys. Rev. A 72, 013603 (2005).
- [Ein25] A. Einstein, *Quantentheorie des einatomigen idealen Gases. Zweite Abhandlung*, Sitzungber. Preuss. Akad. Wiss. **1925**, 3 (1925).
- [Eng06] B. Engeser, A novel surface trapping apparatus for ultracold cesium atoms and the investigation of an Efimov resonance, Ph.D. thesis, Leopold-Franzens-Universität Innsbruck (2006).
- [Esc08] Y. N. Martinez de Escobar, P. G. Mickelson, P. Pellegrini, S. B. Nagel, A. Traverso, M. Yan, R. Cote, and T. C. Killian, *Two-photon photoassociative spectroscopy of ultracold 88-Sr*, arXiv:0808.3434 (2008).
- [Fal04] L. Fallani, L. De Sarlo, J. E. Lye, M. Modugno, R. Saers, C. Fort, and M. Inguscio, Observation of Dynamical Instability for a Bose-Einstein Condensate in a Moving 1D Optical Lattice, Phys. Rev. Lett. 93, 140406 (2004).
- [Fat08a] M. Fattori, C. D'Errico, G. Roati, M. Zaccanti, M. Jona-Lasinio, M. Modugno, M. Inguscio, and G. Modugno, *Atom Interferometry with a Weakly Interacting Bose-Einstein Condensate*, Phys. Rev. Lett. **100**, 080405 (2008).

Bibliography

- [Fat08b] M. Fattori, G. Roati, B. Deissler, C. D'Errico, M. Zaccanti, M. Jona-Lasinio, L. Santos, M. Inguscio, and G. Modugno, *Magnetic Dipolar Interaction in a Bose-Einstein Condensate Atomic Interferometer*, Phys. Rev. Lett. **101**, 190405 (2008).
- [Fer06] G. Ferrari, N. Poli, F. Sorrentino, and G. M. Tino, Long-Lived Bloch Oscillations with Bosonic Sr Atoms and Application to Gravity Measurement at the Micrometer Scale, Phys. Rev. Lett. 97, 060402 (2006).
- [Fer09] F. Ferlaino, S. Knoop, and R. Grimm, *Ultracold Feshbach Molecules*, in: R. V. Krems,
 B. Friedrich, and W. C. Stwalley (Eds.), *Cold Molecules: Theory, Experiment, Applications*, publication expected in March 2009, arXiv:0809.3920.
- [Fes58] H. Feshbach, Unified theory of nuclear reactions, Ann. Phys. 5, 357 (1958).
- [Fli06] A. Flir, *Implementierung und Untersuchung von Raman-Seitenbandkühlung zur Erzeugung eines ultrakalten Cäsiumgases*, Master's thesis, Leopold-Franzens-Universität Innsbruck (2006), available for download at www.ultracold.at.
- [Gab06] G. Gabrielse, D. Hanneke, T. Kinoshita, M. Nio, and B. Odom, *New Determination* of the Fine Structure Constant from the Electron g Value and QED, Phys. Rev. Lett. **97**, 030802 (2006).
- [Gab07] G. Gabrielse, D. Hanneke, T. Kinoshita, M. Nio, and B. Odom, *Erratum: New Determination of the Fine Structure Constant from the Electron g Value and QED*, Phys. Rev. Lett. **99**, 039902 (2007).
- [Ger06] V. Gerginov, K. Calkins, C. E. Tanner, J. J. McFerran, S. Diddams, A. Bartels, and L. Hollberg, *Optical frequency measurements of 6s* ${}^{2}S_{1/2}$ 6p ${}^{2}P_{1/2}$ (D_{1}) transitions in ${}^{133}Cs$ and their impact on the fine-structure constant, Phys. Rev. A **73**, 032504 (2006).
- [Gio02] S. Giovanazzi, A. Görlitz, and T. Pfau, *Tuning the Dipolar Interaction in Quantum Gases*, Phys. Rev. Lett. **89**, 130401 (2002).
- [Gio03] S. Giovanazzi, A. Gorlitz, and T. Pfau, *Ballistic expansion of a dipolar condensate*, J. Opt. B **5**, S208 (2003).
- [Gir60] M. Girardeau, *Relationship between Systems of Impenetrable Bosons and Fermions in One Dimension*, J. Math. Phys. **1**, 516 (1960).
- [GO98] D. Guery-Odelin, J. Soding, P. Desbiolles, and J. Dalibard, *Is Bose-Einstein condensation of atomic cesium possible?*, Europhys. Lett. **44**, 25 (1998).
- [Gre01] M. Greiner, I. Bloch, O. Mandel, T. W. Hänsch, and T. Esslinger, Exploring Phase Coherence in a 2D Lattice of Bose-Einstein Condensates, Phys. Rev. Lett. 87, 160405 (2001).
- [Gre02] M. Greiner, O. Mandel, T. Esslinger, T. W. Hänsch, and I. Bloch, *Quantum phase transition from a superfluid to a Mott insulator in a gas of ultracold atoms*, Nature **415**, 39 (2002).
- [Gri00] R. Grimm, M. Weidemüller, and Y. B. Ovchinnikov, *Optical dipole traps for neutral atoms*, Adv. At. Mol. Opt. Phys. **42**, 95 (2000).

- [Gro61] E. P. Gross, *Structure of a Quantized Vortex in Boson Systems*, Il Nuovo Cimento **20**, 454 (1961).
- [Gup02] S. Gupta, K. Dieckmann, Z. Hadzibabic, and D. E. Pritchard, *Contrast Interferometry using Bose-Einstein Condensates to Measure* h/m and α , Phys. Rev. Lett. **89**, 140401 (2002).
- [Gup03] S. Gupta, *Experiments with degenerate Bose and Fermi gases*, Ph.D. thesis, Massachusetts Institute of Technology (2003).
- [Gus08] M. Gustavsson, E. Haller, M. J. Mark, J. G. Danzl, G. Rojas-Kopeinig, and H.-C. Nägerl, *Control of Interaction-Induced Dephasing of Bloch Oscillations*, Phys. Rev. Lett. 100, 080404 (2008).
- [Ham02] M. Hammes, D. Rychtarik, H.-C. Nägerl, and R. Grimm, *Cold-atom gas at very high densities in an optical surface microtrap*, Phys. Rev. A **66**, 051401 (2002).
- [Han01] D. J. Han, M. T. DePue, and D. S. Weiss, *Loading and compressing Cs atoms in a very far-off-resonant light trap*, Phys. Rev. A **63**, 023405 (2001).
- [Han08] D. Hanneke, S. Fogwell, and G. Gabrielse, *New Measurement of the Electron Magnetic Moment and the Fine Structure Constant*, Phys. Rev. Lett. **100**, 120801 (2008).
- [Har04] T. Hartmann, F. Keck, H. J. Korsch, and S. Mossmann, *Dynamics of Bloch oscillations*, New J. Phys. **6**, 2 (2004).
- [Her03] J. Herbig, T. Kraemer, M. Mark, T. Weber, C. Chin, H.-C. Nägerl, and R. Grimm, *Preparation of a Pure Molecular Quantum Gas*, Science **301**, 1510 (2003).
- [Her05] J. Herbig, *Quantum-Degenerate Cesium: Atoms and Molecules*, Ph.D. thesis, Leopold-Franzens-Universität Innsbruck (2005), available for download at www.ultracold. at.
- [Hop00] S. A. Hopkins, S. Webster, J. Arlt, P. Bance, S. Cornish, O. Maragò, and C. J. Foot, *Measurement of elastic cross section for cold cesium collisions*, Phys. Rev. A 61, 032707 (2000).
- [Hun08] C.-L. Hung, X. Zhang, N. Gemelke, and C. Chin, *Accelerating evaporative cooling of atoms into Bose-Einstein condensation in optical traps*, Phys. Rev. A **78**, 011604 (2008).
- [Ima97] A. Imamoğlu, M. Lewenstein, and L. You, *Inhibition of Coherence in Trapped Bose-Einstein Condensates*, Phys. Rev. Lett. **78**, 2511 (1997).
- [Ino99] S. Inouye, A. P. Chikkatur, D. M. Stamper-Kurn, J. Stenger, D. E. Pritchard, and W. Ketterle, *Superradiant Rayleigh Scattering from a Bose-Einstein Condensate*, Science 285, 571 (1999).
- [Jak98] D. Jaksch, C. Bruder, J. I. Cirac, C. W. Gardiner, and P. Zoller, *Cold Bosonic Atoms in Optical Lattices*, Phys. Rev. Lett. **81**, 3108 (1998).
- [Jak99] D. Jaksch, *Bose-Einstein condensation and applications*, Ph.D. thesis, Leopold-Franzens-Universität Innsbruck (1999).

- [Jav99] J. Javanainen, Phonon approach to an array of traps containing Bose-Einstein condensates, Phys. Rev. A **60**, 4902 (1999).
- [Jul04] P. Julienne, private communication (2004).
- [Kap00] A. E. Kaplan, I. Marzoli, W. E. Lamb, and W. P. Schleich, Multimode interference: Highly regular pattern formation in quantum wave-packet evolution, Phys. Rev. A 61, 032101 (2000).
- [Kas95] A. Kastberg, W. D. Phillips, S. L. Rolston, R. J. C. Spreeuw, and P. S. Jessen, Adiabatic Cooling of Cesium to 700 nK in an Optical Lattice, Phys. Rev. Lett. 74, 1542 (1995).
- [Ker00] A. J. Kerman, V. Vuletić, C. Chin, and S. Chu, Beyond Optical Molasses: 3D Raman Sideband Cooling of Atomic Cesium to High Phase-Space Density, Phys. Rev. Lett. 84, 439 (2000).
- [Ket99] W. Ketterle, D. S. Durfee, and D. M. Stamper-Kurn, Making, probing and understanding Bose-Einstein condensates, in: M. Inguscio, S. Stringari, and C. E. Wieman (Eds.), Proceedings of the International School of Physics - Enrico Fermi, 67, IOS Press, 1999.
- [Koh59] W. Kohn, Analytic Properties of Bloch Waves and Wannier Functions, Phys. Rev. 115, 809 (1959).
- [Köh06] T. Köhler, K. Góral, and P. S. Julienne, *Production of cold molecules via magnetically tunable Feshbach resonances*, Rev. Mod. Phys. **78**, 1311 (2006).
- [Kol02] A. R. Kolovsky, A. V. Ponomarev, and H. J. Korsch, *Damped Bloch oscillations of cold atoms in optical lattices*, Phys. Rev. A **66**, 053405 (2002).
- [Kra02] M. Kraemer, L. Pitaevskii, and S. Stringari, Macroscopic Dynamics of a Trapped Bose-Einstein Condensate in the Presence of 1D and 2D Optical Lattices, Phys. Rev. Lett. 88, 180404 (2002).
- [Kra06] T. Kraemer, M. Mark, P. Waldburger, J. G. Danzl, C. Chin, B. Engeser, A. D. Lange, K. Pilch, A. Jaakkola, H.-C. Nägerl, and R. Grimm, *Evidence for Efimov quantum* states in an ultracold gas of caesium atoms, Nature 440, 315 (2006).
- [Lah07] T. Lahaye, T. Koch, B. Frohlich, M. Fattori, J. Metz, A. Griesmaier, S. Giovanazzi, and T. Pfau, *Strong dipolar effects in a quantum ferrofluid*, Nature **448**, 672 (2007).
- [Leo00] P. J. Leo, C. J. Williams, and P. S. Julienne, *Collision Properties of Ultracold* ¹³³Cs *Atoms*, Phys. Rev. Lett. **85**, 2721 (2000).
- [Li07] W. Li, A. K. Tuchman, H.-C. Chien, and M. A. Kasevich, *Extended Coherence Time with Atom-Number Squeezed States*, Phys. Rev. Lett. **98**, 040402 (2007).
- [Lie03] A. Liem, J. Limpert, H. Zellmer, and A. Tünnermann, 100-W single-frequency master-oscillator fiber power amplifier, Opt. Lett. 28, 1537 (2003).
- [Mad00] K. W. Madison, F. Chevy, W. Wohlleben, and J. Dalibard, *Vortex Formation in a Stirred Bose-Einstein Condensate*, Phys. Rev. Lett. **84**, 806 (2000).

- [Mar07a] M. Mark, Wechselwirkungseffekte eines Cäsium-BEC in eindimensionalen Gittern, Master's thesis, Leopold-Franzens-Universität Innsbruck (2007), available for download at www.ultracold.at.
- [Mar07b] M. Mark, F. Ferlaino, S. Knoop, J. G. Danzl, T. Kraemer, C. Chin, H.-C. Nägerl, and R. Grimm, Spectroscopy of ultracold trapped cesium Feshbach molecules, Phys. Rev. A 76, 042514 (2007).
- [Mar08] M. J. Mark, J. G. Danzl, E. Haller, M. Gustavsson, N. Bouloufa, O. Dulieu, H. Salami, T. Bergeman, H. Ritsch, R. Hart, and H.-C. Nägerl, *Dark resonances for* ground state transfer of molecular quantum gases, arXiv:0811.0695 (2008).
- [Mat99] M. R. Matthews, B. P. Anderson, P. C. Haljan, D. S. Hall, C. E. Wieman, and E. A. Cornell, *Vortices in a Bose-Einstein Condensate*, Phys. Rev. Lett. **83**, 2498 (1999).
- [Men08] C. Menotti, M. Lewenstein, T. Lahaye, and T. Pfau, Dipolar interaction in ultra-cold atomic gases, in: A. Campa, A. Giansanti, G. Morigi, and F. S. Labini (Eds.), Dynamics and thermodynamics of systems with long range interactions: Theory and Experiments, vol. 970 of AIP Conf. Proc., 332–361, AIP, 2008.
- [Met99] H. J. Metcalf and P. van der Straten, *Laser Cooling and Trapping*, Springer Verlag, New York, 1999.
- [Mic05] P. G. Mickelson, Y. N. Martinez, A. D. Saenz, S. B. Nagel, Y. C. Chen, T. C. Killian, P. Pellegrini, and R. Côté, *Spectroscopic Determination of the s-Wave Scattering Lengths of* ⁸⁶Sr and ⁸⁸Sr, Phys. Rev. Lett. **95**, 223002 (2005).
- [Moh08] P. J. Mohr, B. N. Taylor, and D. B. Newell, *CODATA recommended values of the fundamental physical constants: 2006, Rev. Mod. Phys.* **80**, 633 (2008).
- [Mon93] C. R. Monroe, E. A. Cornell, C. A. Sackett, C. J. Myatt, and C. E. Wieman, *Measurement of Cs-Cs elastic scattering at T=30 μK*, Phys. Rev. Lett. **70**, 414 (1993).
- [Mor06] O. Morsch and M. Oberthaler, *Dynamics of Bose-Einstein condensates in optical lattices*, Rev. Mod. Phys. **78**, 179 (2006).
- [Mül05] H. Müller, S. wey Chiow, Q. Long, C. Vo, and S. Chu, *Active sub-Rayleigh alignment* of parallel or antiparallel laser beams, Opt. Lett. **30**, 3323 (2005).
- [Mül08] H. Müller, S. wey Chiow, Q. Long, S. Herrmann, and S. Chu, *Atom Interferometry with up to 24-Photon-Momentum-Transfer Beam Splitters*, Phys. Rev. Lett. **100**, 180405 (2008).
- [Ni08] K.-K. Ni, S. Ospelkaus, M. H. G. de Miranda, A. Pe'er, B. Neyenhuis, J. J. Zirbel, S. Kotochigova, P. S. Julienne, D. S. Jin, and J. Ye, A High Phase-Space-Density Gas of Polar Molecules, Science 322, 231 (2008).
- [Odo06] B. Odom, D. Hanneke, B. D'Urso, and G. Gabrielse, *New Measurement of the Electron Magnetic Moment Using a One-Electron Quantum Cyclotron*, Phys. Rev. Lett. **97**, 030801 (2006).

Bibliography

- [Ped01] P. Pedri, L. Pitaevskii, S. Stringari, C. Fort, S. Burger, F. S. Cataliotti, P. Maddaloni, F. Minardi, and M. Inguscio, *Expansion of a Coherent Array of Bose-Einstein Condensates*, Phys. Rev. Lett. 87, 220401 (2001).
- [Pet00] D. S. Petrov, G. V. Shlyapnikov, and J. T. M. Walraven, *Regimes of Quantum Degeneracy in Trapped 1D Gases*, Phys. Rev. Lett. **85**, 3745 (2000).
- [Pet02] C. J. Pethick and H. Smith, *Bose-Einstein Condensation in Dilute Gases*, Cambridge University Press, 2002.
- [Phi82] W. D. Phillips and H. Metcalf, *Laser Deceleration of an Atomic Beam*, Phys. Rev. Lett. 48, 596 (1982).
- [Pin97] P. W. H. Pinkse, A. Mosk, M. Weidemüller, M. W. Reynolds, T. W. Hijmans, and J. T. M. Walraven, *Adiabatically Changing the Phase-Space Density of a Trapped Bose Gas*, Phys. Rev. Lett. **78**, 990 (1997).
- [Pit61] L. P. Pitaevskii, *Vortex lines in an imperfect Bose gas*, Sov. Phys. JETP **13**, 451 (1961).
- [Pit03] L. Pitaevskii and S. Stringari, *Bose-Einstein Condensation*, Oxford University Press, 2003.
- [Pon06] A. V. Ponomarev and A. R. Kolovsky, *Dipole and Bloch oscillations of cold atoms in a parabolic lattice*, Laser Phys. **16**, 367 (2006).
- [Ree05] J. M. Reeves, O. Garcia, B. Deissler, K. L. Baranowski, K. J. Hughes, and C. A. Sackett, *Time-orbiting potential trap for Bose-Einstein condensate interferometry*, Phys. Rev. A 72, 051605 (2005).
- [Reg03] C. A. Regal, C. Ticknor, J. L. Bohn, and D. S. Jin, *Creation of ultracold molecules from a Fermi gas of atoms*, Nature **424**, 47 (2003).
- [RK07] G. Rojas Kopeinig, A Tunable Bose-Einstein Condensate in a Three-dimensional Optical Lattice Potential, Master's thesis, Leopold-Franzens-Universität Innsbruck (2007), available for download at www.ultracold.at.
- [Roa04] G. Roati, E. de Mirandes, F. Ferlaino, H. Ott, G. Modugno, and M. Inguscio, Atom Interferometry with Trapped Fermi Gases, Phys. Rev. Lett. 92, 230402 (2004).
- [Ryc04] D. Rychtarik, B. Engeser, H.-C. Nägerl, and R. Grimm, *Two-Dimensional Bose-Einstein Condensate in an Optical Surface Trap*, Phys. Rev. Lett. **92**, 173003 (2004).
- [Sag05] J. M. Sage, S. Sainis, T. Bergeman, and D. DeMille, *Optical Production of Ultracold Polar Molecules*, Phys. Rev. Lett. **94**, 203001 (2005).
- [Sal91] B. E. A. Saleh and M. C. Teich, *Fundamentals of photonics*, Wiley, 1991.
- [Sme03] A. Smerzi and A. Trombettoni, *Nonlinear tight-binding approximation for Bose-Einstein condensates in a lattice*, Phys. Rev. A **68**, 023613 (2003).
- [Söd98] J. Söding, D. Guéry-Odelin, P. Desbiolles, G. Ferrari, and J. Dalibard, *Giant Spin Relaxation of an Ultracold Cesium Gas*, Phys. Rev. Lett. **80**, 1869 (1998).
- [Ste02] D. A. Steck, *Cesium D Line Data*, http://steck.us/alkalidata , revision 2.1 (2002), and references therein.
- [Tha06] G. Thalhammer, K. Winkler, F. Lang, S. Schmid, R. Grimm, and J. H. Denschlag, Long-Lived Feshbach Molecules in a Three-Dimensional Optical Lattice, Phys. Rev. Lett. 96, 050402 (2006).
- [Tho03] A. M. Thomas, S. Hopkins, S. L. Cornish, and C. J. Foot, Strong evaporative cooling towards Bose–Einstein condensation of a magnetically trapped caesium gas, J. Opt. B 5, S107 (2003).
- [Tie92] E. Tiesinga, A. J. Moerdijk, B. J. Verhaar, and H. T. C. Stoof, Conditions for Bose-Einstein condensation in magnetically trapped atomic cesium, Phys. Rev. A 46, R1167 (1992).
- [Tie93] E. Tiesinga, B. J. Verhaar, and H. T. C. Stoof, *Threshold and resonance phenomena in ultracold ground-state collisions*, Phys. Rev. A **47**, 4114 (1993).
- [Tre01] P. Treutlein, K. Y. Chung, and S. Chu, *High-brightness atom source for atomic fountains*, Phys. Rev. A **63**, 051401 (2001).
- [Ude00] T. Udem, J. Reichert, T. W. Hänsch, and M. Kourogi, *Absolute optical frequency measurement of the cesium D2 line*, Phys. Rev. A **62**, 031801 (2000).
- [Unt05] P. Unterwaditzer, Aufbau eines vollständigen Diodenlasersystems zur Laserkühlung und Detektion von gespeicherten Cs-Atomen, Master's thesis, Leopold-Franzens-Universität Innsbruck (2005), available for download at www.ultracold.at.
- [Vol06] T. Volz, N. Syassen, D. Bauer, E. Hansis, S. Dürr, and G. Rempe, Preparation of a quantum state with one molecule at each site of an optical lattice, Nat. Phys. 2, 692 (2006).
- [Vul99] V. Vuletić, A. J. Kerman, C. Chin, and S. Chu, Observation of Low-Field Feshbach Resonances in Collisions of Cesium Atoms, Phys. Rev. Lett. 82, 1406 (1999).
- [Wap03] A. H. Wapstra, G. Audi, and C. Thibault, *The 2003 atomic mass evaluation:* (I). Evaluation of input data, adjustment procedures, Nucl. Phys. A **729**, 129 (2003).
- [Was93] C. Waschke, H. G. Roskos, R. Schwedler, K. Leo, H. Kurz, and K. Köhler, Coherent submillimeter-wave emission from Bloch oscillations in a semiconductor superlattice, Phys. Rev. Lett. 70, 3319 (1993).
- [Web03a] T. Weber, *Bose-Einstein Condensation of Optically Trapped Cesium*, Ph.D. thesis, Leopold-Franzens-Universität Innsbruck (2003), available for download at www. ultracold.at.
- [Web03b] T. Weber, J. Herbig, M. Mark, H.-C. Nägerl, and R. Grimm, *Bose-Einstein Condensation of Cesium*, Science **299**, 232 (2003).
- [Web03c] T. Weber, J. Herbig, M. Mark, H.-C. Nägerl, and R. Grimm, *Three-Body Recombination at Large Scattering Lengths in an Ultracold Atomic Gas*, Phys. Rev. Lett. 91, 123201 (2003).

- [Wic02] A. Wicht, J. M. Hensley, E. Sarajlic, and S. Chu, *A Preliminary Measurement of the Fine Structure Constant Based on Atom Interferometry*, Phys. Scr. **T102**, 82 (2002).
- [Win07a] K. Winkler, Ultracold Molecules and Atom Pairs in Optical Lattice Potentials, Ph.D. thesis, Leopold-Franzens-Universität Innsbruck (2007).
- [Win07b] K. Winkler, F. Lang, G. Thalhammer, P. v. d. Straten, R. Grimm, and J. H. Denschlag, Coherent Optical Transfer of Feshbach Molecules to a Lower Vibrational State, Phys. Rev. Lett. 98, 043201 (2007).
- [Wit05] D. Witthaut, M. Werder, S. Mossmann, and H. J. Korsch, *Bloch oscillations of Bose-Einstein condensates: Breakdown and revival*, Phys. Rev. E **71**, 036625 (2005).
- [Wu01] B. Wu and Q. Niu, *Landau and dynamical instabilities of the superflow of Bose-Einstein condensates in optical lattices*, Phys. Rev. A **64**, 061603 (2001).
- [Wu03] B. Wu and Q. Niu, Superfluidity of Bose-Einstein condensate in an optical lattice: Landau-Zener tunnelling and dynamical instability, New J. Phys. 5, 104 (2003).
- [Zhe04] Y. Zheng, M. Koštrun, and J. Javanainen, *Low-Acceleration Instability of a Bose-Einstein Condensate in an Optical Lattice*, Phys. Rev. Lett. **93**, 230401 (2004).

A

Publications

Control of Interaction-Induced Dephasing of Bloch Oscillations

M. Gustavsson, E. Haller, M. J. Mark, J. G. Danzl, G. Rojas-Kopeinig, and H.-C. Nägerl

Institut für Experimentalphysik und Forschungszentrum für Quantenphysik, Universität Innsbruck, 6020 Innsbruck, Austria

(Received 26 October 2007; published 28 February 2008)

We report on the control of interaction-induced dephasing of Bloch oscillations for an atomic Bose-Einstein condensate in an optical lattice. We quantify the dephasing in terms of the width of the quasimomentum distribution and measure its dependence on time for different interaction strengths which we control by means of a Feshbach resonance. For minimal interaction, the dephasing time is increased from a few to more than 20 thousand Bloch oscillation periods, allowing us to realize a BECbased atom interferometer in the noninteracting limit.

DOI: 10.1103/PhysRevLett.100.080404

PACS numbers: 03.75.Dg, 05.30.Jp, 34.50.-s, 37.10.Jk

Ultracold atomic systems have initiated a revolution in the field of precision measurements. Laser cooled thermal samples are used for ultrahigh resolution laser spectroscopy [1], are at the heart of modern atomic fountain clocks [2,3], and allow for the realization of matter-wave interferometers for high-precision inertial sensing [4] and highprecision determination of fundamental constants [5]. Atomic Bose-Einstein condensates (BEC), the matterwave analog to the laser, combine high brightness with narrow spatial and momentum spread. In general, the resolution is limited only by the quantum mechanical uncertainty principle, and BECs could thus serve as ideal sources for precision measurements and, in particular, for matter-wave interferometers [6]. Atom-atom interactions, however, have to be taken into account, as they lead to collisional dephasing and give rise to density dependent mean-field shifts in the interferometric signal. It is thus advisable to either operate a BEC-based atom interferometer in the dilute density limit, possibly sacrificing a high signal-to-noise ratio, or to find ways of reducing or even nulling the strength of the interaction altogether. Precisely the latter is feasible in the vicinity of magnetically induced Feshbach resonances where the atomic s-wave scattering length and hence the strength of the atom-atom contact interaction go through a zero crossing [7]. It is thus possible to experimentally investigate the reduction and even disappearance of interaction-induced effects on the interferometric signal as the scattering length is tuned towards zero by means of an externally controlled magnetic field.

A paradigm atom interferometric effect is the wellknown phenomenon of Bloch oscillations [8]. Bloch oscillations for the mean quasimomentum are the result of single atom interference as the atomic wave packet, subject to a constant force, is Bragg reflected in the presence of a periodic optical lattice potential. They have been observed for ultracold thermal samples [5,8-10], for atoms in interacting BECs [11,12], and for ensembles of noninteracting quantum-degenerate fermions [12]. For the case of the interacting BEC, strong dephasing is found as evidenced by a rapid broadening and apparent smearing out of the momentum distribution in the first Brillouin zone, limiting the observation of Bloch oscillations to a few cycles for typical atomic densities in a BEC. In addition, the measured initial width of the momentum distribution is comparable to the extent of the Brillouin zone, as interaction energy is converted into kinetic energy upon release of the BEC from the lattice potential, thus greatly reducing the contrast of the oscillations [12].

In this Letter, we report on the control of interactioninduced dephasing of Bloch oscillations for a BEC in a vertically oriented optical lattice under the influence of gravity. Control is obtained by means of a zero -crossing for the atomic s-wave scattering length a. We observe the transition from an interacting BEC to a noninteracting BEC by measuring the rate of dephasing, given by the change of the width of the momentum distribution, as a function of a. We identify a clear minimum for the dephasing which we associate with the zero crossing for a. At the minimum, more than 2×10^4 oscillations can be observed with high contrast, and the zero crossing can be determined with high precision. For our measurements at nonzero scattering length, we greatly reduce broadening of the momentum distribution by rapidly switching the interaction strength to zero upon release from the lattice potential. Our measurements indicate that BECs can indeed be used as a source for precision atom interferometry, as effects of the interaction can be greatly reduced. For a noninteracting BEC, we intentionally induce dephasing by means of a weak optical force gradient and observe collapse and revivals of Bloch oscillations.

The starting point for our experiments is an essentially pure BEC with typically 1×10^5 Cs atoms in the $|F = 3, m_F = 3\rangle$ hyperfine ground state sublevel confined in a crossed-beam dipole trap generated by one vertically $(L_1, \text{ with } 1/e^2$ -beam diameter 256 μ m) and one more tightly focused horizontally $(L_2, \text{ with diameter 84 } \mu\text{m})$ propagating laser beam at a wavelength near 1064 nm. We support the optical trapping by magnetic levitation against gravity [13]. For BEC preparation, we basically follow the procedure described in Refs. [13,14]. The

0031-9007/08/100(8)/080404(4)

strength of the interaction can be tuned by means of a broad Feshbach resonance with a pole at -11.7 G. The resonance causes a zero crossing for the scattering length a near an offset magnetic field value of 17 G with a slope of 61 a0/G [15]. Here, a_0 denotes Bohr's radius. The lattice potential is generated by a vertically oriented standing laser wave generated by retro-reflection, collinear with L_1 , but with much larger diameter of 580 μ m. This allows independent control of lattice depth and radial (i.e., horizontal) confinement. The light comes from a home-built single-mode fiber amplifier [16] seeded with highly-stable light at $\lambda = 1064.4946(1)$ nm. We turn on the optical lattice potential exponentially to a depth of $7.9E_R$ within 1000 ms, where $E_R = h^2/(2m\lambda^2) = k_B \times 64$ nK is the photon recoil energy and m is the mass of the Cs atom. The slow ramp assures that the BEC is adiabatically loaded into the lowest Bloch band of the lattice, and it avoids horizontal excitations. We load between 40 to 65 lattice sites, depending on the initial vertical extent of the BEC. We then reduce the power in L_2 to zero within 300 μ s. Subsequently, the magnetic field gradient needed for levitation is ramped down, and a bias magnetic field is tuned to the desired value within 100 μ s. For the present experiments, we adjust a in the range from -2 to $300 a_0$ with magnetic bias fields from 17 to 23 G. The step in a leads to some unavoidable horizontal excitation as a result of the change of the Thomas-Fermi profile. We control the average bias field to about 1 mG. The confinement of the BEC in the lattice as given by L_1 gives horizontal trapping frequencies in the range of 5 to 10 Hz. We then let the atoms evolve in the lattice under the influence of the gravitational force for variable hold time T. Finally, we switch off the horizontal confinement and ramp the lattice depth adiabatically to zero within 300 μ s to measure the momentum distribution by the standard time-of-flight technique, taking an absorption picture on a CCD camera. For some of the data, we turn on the magnetic levitation field to allow for longer expansion times up to 100 ms. To minimize broadening of the distribution as a result of interaction, we switch the scattering length to zero during the release and the initial time-of-flight.

We observe persistent Bloch oscillations when minimizing the effect of interactions at a magnetic field value of 17.12 G (see below). Figures 1(a)-1(d) show the evolution of the momentum distribution during the first, the 1000th, the 10 000th, and the 20 000th Bloch cycle. Initially, the momentum distribution exhibits narrow peaks. Their full width Δp [17] is as narrow as about 0.15 $\hbar k$, where $k = 2\pi/\lambda$. Very little broadening along the vertical direction is seen after the first 1000 cycles. Initial excitation of horizontal motion as a result of ramping the power in L_2 and switching the scattering length leads to some horizontal spreading. After 20 000 cycles, the distribution has started to spread out noticeably along the vertical direction.



FIG. 1. Long-lived Bloch oscillations for a noninteracting BEC with Cs atoms in the vertical lattice under the influence of gravity. Each picture shows one Bloch cycle in successive time-of-flight absorption images corresponding to the momentum distribution at the time of release from the lattice. Displayed are the first (a), the 1000th (b), the 10 000th (c), and the 20 000th (d) Bloch cycle for minimal interaction near the zero crossing for the scattering length.

Figure 2 highlights the high number of Bloch oscillations, which we can observe for the case of minimal interaction strength. It shows how the strongest peak of the momentum distribution cycles through the first Brillouin zone with the typical sawtooth behavior [8]. More than 20 000 cycles can easily be followed. From a fit to the data, we determine the Bloch period to 0.575 380 7(5) ms. Assuming that no additional forces act on the sample, the local gravitational constant is g =9.803 821(9) m/s². The error is statistical only. While we took care to minimize magnetic field gradients, we expect them to be the dominant contribution to the systematic error.

In order to quantify the dephasing of Bloch oscillations, we determine for each Bloch period the width Δp of the momentum distribution at the instant in time when the peak of the distribution is centered at zero momentum, i.e., for the central picture of each series shown in Fig. 1. Figure 3(a) displays Δp up to the 300th Bloch cycle for different interaction strengths ranging from 0 to 300 a_0 . For minimal interaction strength ($a \approx 0 a_0$), we see no broadening of the distribution. Broadening can clearly be



FIG. 2 (color online). Position of the strongest peak in the momentum distribution as a function of the number N of Bloch oscillations (dots). More than 20 000 cycles can be followed with high contrast. A fit to the data (solid curve) yields a Bloch period of 0.575 3807(5) ms.



FIG. 3 (color online). Width Δp of the momentum distribution for different interaction strengths. (a) Evolution of Δp as a function of the number N of Bloch cycles for different values of the scattering length ($a = 0, 25, 50, 100, \text{ and } 300a_0$ from bottom (full circles) to top (open squares). The solid curves are derived from a numerical model calculation, see text. (b) Width Δp for a fixed number of cycles N = 1 (full circles), 25 (full squares), 50 (full diamonds), 100 (open circles), 150 (open squares), and 200 (open diamonds) as a function of scattering length. The solid line represents the model calculation. All error bars correspond to ± 1 standard deviation resulting from 7 measurements. The data and the simulations correspond to the following parameters: lattice depth: $7.9E_R$, scattering length during lattice loading: 210 a_0 , trapping frequencies in L_1 and L_2 : 10 and 8 Hz, atom number in the BEC: 5×10^4 .

seen for $a = 25 a_0$, and the rate of broadening then increases with increasing interaction strength. For $a \ge a$ 50 a_0 , the width Δp saturates within the chosen observation time to a value of about $1.3\hbar k$ as the momentum distribution completely fills the first Brillouin zone [18]. To a good approximation, we find that Δp initially increases linearly with time. In Fig. 3(b), we plot Δp as a function of interaction strength for various fixed numbers of Bloch cycles. Δp appears to scale with the square root of the interaction strength. Both observations agree well with a simple model for the dephasing of Bloch oscillations, which predicts $\Delta p \propto \sqrt{aT}$ [19] for sufficiently short times T. In order to verify this model, we have performed numerical calculations solving the one-dimensional Gross-Pitaevskii equation in the presence of an optical lattice under the influence of gravity for the typical parameters of our experiment according to the method detailed in Ref. [20]. Via Fourier transform of the spatial wave function, we determine the momentum distribution and its width. As shown in Fig. 3 (solid lines), we find very good agreement with our measurements with no adjustable parameters when we add a constant offset of $0.1\hbar k$ to all

the numerical curves. This offset takes into account residual interactions during release from the lattice as a result of the finite magnetic switching speed, which leads to some artificial broadening of the distribution. We attribute the systematic discrepancy for the N = 50 data in Fig. 3(b) to the horizontal motion which leads to modulations in the density that adds a modulation onto Δp also seen in Fig. 3(a).

To find the value for the magnetic field that gives minimal broadening, we measure Δp after 6951 cycles in the vicinity of the crossing. Figure 4 plots Δp as a function of magnetic field. It shows a clear minimum, which we expect to correspond to the zero crossing for the scattering length. From a Gaussian fit, we determine the center position of the minimum to be at 17.119(2) G. The one-sigma error takes into account our statistical error in magnetic field calibration. To our knowledge, this is the most precise determination of a minimum for the elastic cross section in ultracold atom scattering. We believe that our measurements are limited by the ambient magnetic field noise, leading to a finite width for the distribution of the scattering length. In fact, a reduction of the atomic density gives longer decay times for the Bloch oscillations. Note that in the scattering length regime considered here, the effect of the (magnetic) dipole-dipole interaction [21] should start to play a role.

Our capability to observe Bloch oscillations on extended time scales without interaction-induced dephasing allows us to study the effect of deliberately imposed dephasing. For this, we apply a linear force gradient ∇F corresponding to harmonic trapping at $\nu = 40(1)$ Hz along the vertical direction by turning on laser beam L_2 during the hold time. Figure 5 shows the widths Δp for two cycle phases separated by π as a function of the number N of Bloch cycles. The two phases correspond to the single- resp.



FIG. 4. Broadening of the momentum distribution as a result of 6951 Bloch oscillations near the zero crossing for the scattering length. The width Δp is plotted as a function of magnetic field (dots). The solid line is a Gaussian fit with a rms-width of 4.5 mG. The fit is centered at 17.119(2) G. The zero for the scattering length scale on top was chosen to agree with this value.



FIG. 5. Collapse and revival of Bloch oscillations for the case of a noninteracting BEC with a vertical force gradient. The width Δp is plotted as a function of the number *N* of Bloch cycles for two cycle phases separated by π . For selected cycles (N = 1, 70, 140, 210, and 280), two absorption images corresponding to the two cycle phases are shown.

symmetric double-peaked distribution. Both widths rapidly increase resp. decrease to the same value of $1.3\hbar k$ within about N = 30 oscillations. Here the ensemble is dephased. It then remains dephased for about 200 cycles. Partial rephasing at intermediate times not reflected in the widths can be seen from the absorption images. Revival of the oscillations [22] happens around N = 280 when the values for both widths separate again [23]. This number agrees well with the expected value of $N_{\rm rev} = 292(15)$ given by $N_{\rm rev} = F_{\rm grav}/(\nabla F d) = mg/(m\omega^2 d)$, where $F_{\rm grav}$ is the gravitational force, $\omega = 2\pi\nu$, and $d = \lambda/2$ is the lattice spacing. Subsequently, the widths collapse again to the common value. In further measurements, we see up to four collapses and revivals.

In summary, we have demonstrated the control of interaction-induced dephasing near a zero crossing for the scattering length. On the crossing, we have realized a noninteracting BEC, which allows us to observe more than 20000 Bloch cycles, indicating a matter-wave coherence time of more than 10 s. The broadening of the momentum distribution agrees well with results from theoretical models. We believe that the number of observable Bloch cycles is limited by residual interactions as a result of magnetic field noise. Our results open up exciting new avenues for precision measurements with quantum-degenerate gases. For example, it is now possible to perform sensitive measurements of forces on short length scales, such as the Casimir-Polder force near a dielectric surface [24]. Future experimental work can now address the nature of the dephasing [25] by studying structure in the momentum distribution.

A similar experiment on long-lasting Bloch oscillations and control of the interaction strength has recently been performed with a BEC of ³⁹K atoms at LENS, Italy [26]. We thank A. Daley for theoretical support and for help with setting up the numerical calculations and A. Buchleitner and his group for useful discussions. We are grateful to A. Liem and H. Zellmer for valuable assistance in setting up the 1064 nm fiber amplifier system. We acknowledge contributions by P. Unterwaditzer and T. Flir during the early stages of the experiment. We are indebted to R. Grimm for generous support and gratefully acknowledge funding by the Austrian Ministry of Science and Research (BMWF) and the Austrian Science Fund (FWF).

- [1] S. A. Diddams et al., Science 306, 1318 (2004).
- [2] S. Bize et al., J. Phys. B 38, S449 (2005).
- [3] M. M. Boyd et al., Phys. Rev. Lett. 98, 083002 (2007).
- [4] A. Peters, K. Y. Chung, and S. Chu, Nature (London) 400, 849 (1999).
- [5] P. Cladé et al., Phys. Rev. Lett. 96, 033001 (2006).
- [6] S. Gupta et al., Phys. Rev. Lett. 89, 140401 (2002).
- [7] For a review, see T. Köhler, K. Góral, and P.S. Julienne, Rev. Mod. Phys. 78, 1311 (2006).
- [8] M. Ben Dahan et al., Phys. Rev. Lett. 76, 4508 (1996).
- [9] R. Battesti et al., Phys. Rev. Lett. 92, 253001 (2004).
- [10] G. Ferrari et al., Phys. Rev. Lett. 97, 060402 (2006).
- [11] O. Morsch et al., Phys. Rev. Lett. 87, 140402 (2001).
- [12] G. Roati et al., Phys. Rev. Lett. 92, 230402 (2004).
- [13] T. Weber et al., Science 299, 232 (2003).
- [14] T. Kraemer et al., Appl. Phys. B 79, 1013 (2004).
- [15] P. Julienne (private communication).
- [16] A. Liem et al., Opt. Lett. 28, 1537 (2003).
- [17] We define the full width Δp to be the root-mean-square (rms) diameter of the distribution.
- [18] The momentum distribution of a fully dephased interacting ensemble carries high-contrast substructure which will be discussed in a forthcoming publication.
- [19] D. Witthaut *et al.*, Phys. Rev. E **71**, 036625 (2005). Note that for the quasi-one-dimensional case, the interaction constant in Eq. (33) is proportional to \sqrt{a} , see [20].
- [20] A. Smerzi and A. Trombettoni, Phys. Rev. A 68, 023613 (2003).
- [21] S. Giovanazzi, A. Görlitz, and T. Pfau, Phys. Rev. Lett. 89, 130401 (2002).
- [22] A. V. Ponomarev and A. R. Kolovsky, Laser Phys. 16, 367 (2006).
- [23] Note that the revived Bloch cycles are subject to a phase shift, which depends on the vertical location of the harmonic trap minimum with respect to the lattice minima.
- [24] I. Carusotto et al., Phys. Rev. Lett. 95, 093202 (2005).
- [25] A. Buchleitner and A. R. Kolovsky, Phys. Rev. Lett. 91, 253002 (2003).
- [26] M. Fattori et al., Phys. Rev. Lett. 100, 080405 (2008).

REPORTS



Fig. 3. Mass function of the stars formed in the simulation illustrated in Fig. 1. The stars form with masses close to $0.1 M_{\odot}$ but grow quickly through gas accretion. The mass function therefore has a peak at ~0.8 M_{\odot} , above which it has a power-law form with a slope comparable to that of the Salpeter slope, illustrated by the diagonal line.

involve further star formation or self-gravity–driven accretion. Assuming that all the gas would accrete directly onto the black hole on a viscous time scale (21) on the order of 10^7 years yields an average mass accretion rate on the order of $10^{-4} M_{\odot}$ /year for the $10^4 M_{\odot}$ cloud and $10^{-3} M_{\odot}$ /year for the $10^5 M_{\odot}$ cloud. This gives a maximum accretion luminosity on the order of 10^{43} ergs/s, or about 1% of the Eddington luminosity. This additional source of radiation, and that from the newly formed stars, could increase the gas temperature in the disk and thus increase the fragment masses.

Our simulations show that an infalling molecular cloud can indeed form an eccentric disk around a supermassive black hole, and that although the tidal force of the black hole will dis-



Fig. 4. Mass function of the stars formed in the simulation illustrated in Fig. 2. The mass function is extremely top-heavy and appears to have two populations of stars: one population of massive stars with masses from 10 M_{\odot} to 100 M_{\odot} , and another with masses between 1 M_{\odot} and 10 M_{\odot} .

rupt the cloud, it does not destroy the small-scale structures that seed the disk fragmentation. Furthermore, the compressional heating of the infalling gas results in the formation of a population of stars biased toward higher masses. The stellar masses depend crucially on the mass of the infalling cloud and on its impact parameter, allowing for a variety of final outcomes. The initial disk eccentricity also means that the stars can form with initial eccentricities and, if the molecular cloud is sufficiently massive, the stars that form may be extremely massive. This is therefore a viable mechanism for forming the rings of young, massive stars within ~0.1 pc of the galactic center. What is still unclear, however, is the origin of the infalling cloud and the probability of the small impact parameter that is required.

References and Notes

- 1. R. Genzel et al., Mon. Not. R. Astron. Soc. 291, 219 (1997).
- 2. R. Schödel et al., Nature 419, 694 (2002).
- 3. A. Ghez et al., Astrophys. J. 620, 744 (2005).
- 4. T. Paumard *et al., Astrophys. J.* **643**, 1011 (2006).
- 5. S. Nayakshin, R. Sunyaev, *Mon. Not. R. Astron. Soc.* **364**, L23 (2005).
- 6. B. Elmegreen, Astrophys. J. 648, 572 (2006).
- E. S. Phinney, in *The Center of the Galaxy: Proceedings of* the 136th Symposium of the IAU, M. Morris, Ed. (Kluwer, Dordrecht, Netherlands, 1989), p. 543–553.
- 8. Y. Levin, A. M. Beloborodov, *Astrophys. J.* **590**, L33 (2003).
- S. Nayakshin, J. Cuadra, Astron. Astrophys. 437, 437 (2005).
- M. A. Gürkan, F. A. Rasio, Astrophys. J. 628, 236 (2005).
 S. Nayakshin, J. Cuadra, V. Springel, Mon. Not. R. Astron.
- Soc. **379**, 21 (2007). 12. C. F. Gammie, Astrophys. J. **553**, 174 (2001).
- 13. W. K. M. Rice *et al.*, *Mon. Not. R. Astron. Soc.* **339**, 1025
- (2003). 14. R. D. Alexander, M. C. Begelman, P. J. Armitage,
- Astrophys. J. 654, 907 (2007).
- 15. R. D. Alexander, P. J. Armitage, J. Cuadra,
- M. C. Begelman, Astrophys. J. 674, 927 (2008).
 M. Wardle, F. Yusef-Zadeh, Astrophys. J. 683, L37 (2008).
- I. Monaghan, Annu. Rev. Astron. Astrophys. 30, 543 (1992).
- D. Stamatellos, A. P. Whitworth, T. Bisbas, S. Goodwin, Astron. Astrophys. 475, 37 (2007).
- M. R. Bate, I. A. Bonnell, N. M. Price, *Mon. Not. R.* Astron. Soc. 277, 362 (1995).
- See supporting material on *Science* Online.
 The viscous time scale is the time for shearing
- interactions between adjacent annuli to transport mass to the central black hole.

Supporting Online Material

www.sciencemag.org/cgi/content/full/321/5892/1060/DC1 Materials and Methods Figs. S1 to S3 References

19 May 2008; accepted 24 July 2008 10.1126/science.1160653

Quantum Gas of Deeply Bound Ground State Molecules

Johann G. Danzl,¹* Elmar Haller,¹ Mattias Gustavsson,¹ Manfred J. Mark,¹ Russell Hart,¹ Nadia Bouloufa,² Olivier Dulieu,² Helmut Ritsch,³ Hanns-Christoph Nägerl¹

Molecular cooling techniques face the hurdle of dissipating translational as well as internal energy in the presence of a rich electronic, vibrational, and rotational energy spectrum. In our experiment, we create a translationally ultracold, dense quantum gas of molecules bound by more than 1000 wave numbers in the electronic ground state. Specifically, we stimulate with 80% efficiency, a two-photon transfer of molecules associated on a Feshbach resonance from a Bose-Einstein condensate of cesium atoms. In the process, the initial loose, long-range electrostatic bond of the Feshbach molecule is coherently transformed into a tight chemical bond. We demonstrate coherence of the transfer in a Ramsey-type experiment and show that the molecular sample is not heated during the transfer. Our results show that the preparation of a quantum gas of molecules in specific rovibrational states is possible and that the creation of a Bose-Einstein condensate of molecules in their rovibronic ground state is within reach.

Using the samples of molecules are ideally suited for fundamental studies in physics and chemistry, ranging from few-body collisional physics (1-4), ultracold chemistry

(5), and high-resolution spectroscopy (6, 7) to quantum gas preparation, molecular Bose-Einstein condensation (δ), and quantum processing (9). For many of the proposed experiments, full con-

trol over the molecular wave function in specific deeply bound rovibrational states is needed. High densities are required for molecular quantum gas studies. Only in the rovibronic ground state (the lowest vibrational and rotational energy level of the electronic ground state) is collisional stability assured. However, direct molecular cooling toward high phase-space densities seems yet out of reach (10), whereas techniques such as Feshbach association (11) and photoassociation (12) either produce molecules exclusively in weakly bound rovibrational levels or suffer from low production rates and low state selectivity.

To produce a quantum gas of molecules in their absolute ground state, Jaksch *et al.* (13) proposed a scheme for homonuclear alkali molecules in which the technique of stimulated

¹Institut für Experimentalphysik und Zentrum für Quantenphysik, Universität Innsbruck, Technikerstraße 25, 6020 Innsbruck, Austria. ²Laboratoire Aimé Cotton, CNRS, Université Paris-Sud Bâtiment 505, 91405 Orsay Cedex, France. ³Institut für Theoretische Physik und Zentrum für Quantenphysik, Universität Innsbruck, Technikerstraße 25, 6020 Innsbruck, Austria.

^{*}To whom correspondence should be addressed. E-mail: johann.danzl@uibk.ac.at

two-photon transfer is repeatedly applied to molecules associated from a high-density sample of ultracold atoms. The initially very loosely bound molecules are transferred in successive steps to the rovibrational ground state of the singlet $X^1 \Sigma_{\sigma}^+$ molecular potential. The advantages of this scheme are that it is fully coherent, not relying on spontaneous processes, and that it involves only a very small number of intermediate levels. It promises that a ground state binding energy typically of 0.5 eV can be carried away without heating the molecular sample. It essentially preserves phase-space density, allowing the molecular sample to inherit the high initial phase-space density from the atomic sample. However, to realize this scheme, several challenges have to be met. First, there is a large difference in internuclear separation that has to be bridged: The overlap between the radial wave function of the least bound molecules and the radial wave functions of

deeply bound molecular levels is extremely low, potentially leading to prohibitively low transition rates for the two-photon transitions. Second, the scheme requires the identification of suitable intermediate molecular levels while strictly avoiding parasitic excitations. Third, a large difference in binding energy has to be overcome. On a more technical side, the lasers driving the two-photon transitions at widely different wavelengths need to have extremely low relative short-term phase jitter and high longterm frequency stability to allow for coherence and reproducibility. In important experiments, Winkler et al. (14) and, recently, Ospelkaus et al. (15) demonstrated highly efficient two-photon transfer into lower-lying molecular levels starting from weakly bound dimer molecules, which were associated from ultracold atoms on a Feshbach resonance (11). However, the transferred molecules are still weakly bound. Their binding energy, on the order of the atomic hy-



Fig. 1. (**A**) Molecular level scheme for Cs₂. Molecules in a weakly bound Feshbach level are transferred to rovibrational level |v = 73, J = 2> of the singlet $X^{1}\Sigma_{g}^{+}$ potential with a binding energy of 1061 cm⁻¹ in a two-photon STIRAP process with wavelengths near 1126 and 1006 nm via the 225th level of the electronically excited $(A^{1}\Sigma_{u}^{+} - b^{3}\Pi_{u})0_{u}^{+}$ potentials. The $X^{1}\Sigma_{g}^{+}$ potential has about 155 vibrational levels. a_{0} is the Bohr radius. (**B**) Zeeman diagram showing the energy of all relevant weakly bound molecular levels for initial Feshbach molecular state preparation (*18*). The binding energy is given with respect to the F = 3, $m_{\rm F} = 3$ two-atom asymptote. The molecules are produced on a *d*-wave Feshbach resonance at 4.8 mT (inset) and then transferred to the weakly bound *s*-wave state |s> on an avoided state crossing. Further lowering of the magnetic offset field to 1.9 mT transfers the molecules from state |s> to state |a>, the starting state for the STIRAP transfer. (**C**) STIRAP transfer scheme (*19*). The molecules are transferred from the initial state |a> to the final state |g> = |v = 73, J = 2> by means of two overlapping laser pulses for which laser L_{2} is pulsed on before L_{1} . The detunings and Rabi frequencies of L_{i} are Δ_{i} and Ω_{i} , i = 1, 2.

perfine splitting, is $<10^{-4}$ of the binding energy of the rovibrational ground state, and wave function overlap with this state is still negligible.

In this experiment, we demonstrate the crucial step toward full control of the molecular wave function and toward the formation of a Bose-Einstein condensate (BEC) of molecules in their rovibronic ground state by linking weakly bound molecular states with deeply bound rovibrational states. We coherently transfer an ultracold quantum gas of weakly bound cesium Feshbach molecules to the rovibrational level |v = 73, J = 2 of the singlet $X^1 \Sigma_{\alpha}^+$ potential, bound by 1061 cm⁻¹ (or $h \times 31.81$ THz, where h is Planck's constant), corresponding to more than one-fourth of the binding energy of the rovibrational ground state. To achieve this result, we overcome low wave function overlap by using a suitable intermediate excited molecular state while avoiding excitation into loss channels, and we reference the transfer lasers to a frequency comb, allowing us to flexibly bridge bindingenergy differences of more than 1000 cm^{-1} .

Figure 1 shows the energy of the relevant molecular and atomic states. Our experiment starts with a cigar-shaped BEC of cesium atoms in the lowest hyperfine sublevel F = 3, $m_{\rm F} = 3$ in an optical dipole trap. For BEC production, we essentially follow the procedure detailed in (16). For Feshbach molecule production out of the BEC, we ramp up the offset magnetic field from the initial value of 2.1 mT to ~5.0 mT in 10 ms. We then ramp down, sweeping across a d-wave Feshbach resonance at 4.8 mT after ~1 ms, as shown in Fig. 1B (17, 18). Our procedure [see (17)] gives an ultracold and dense sample of up to 11,000 molecules every 10 s at densities above 1×10^{11} cm⁻³. For the statetransfer experiments discussed here, we do not separate the molecules from the original BEC. Upon lowering the magnetic field, the molecules are transferred from the initial state $|d\rangle$ to a still weakly bound *s*-wave molecular state $|s\rangle$ of the lowest hyperfine channel ($F_1 = 3, F_2 = 3$) via an avoided crossing (18). The index i = 1, 2denotes the *i*th atom.

Upon further lowering the magnetic field to about 2.2 mT, the molecules enter into a closed channel s-wave molecular state $|a\rangle$ via a second, broad avoided crossing (18). This state belongs to the uppermost hyperfine channel ($F_1 =$ 4, $F_2 = 4$) and thus has an effective binding energy of more than $2 \times hv_{Cs}$. Here, $v_{Cs} \approx 9.19$ GHz is the Cs clock frequency. Similar to $|s\rangle$, this state is a mixture of the $X^1 \Sigma_{\rho}^+$ ground state and the lowest triplet $a^{3}\Sigma_{u}^{+}$ state, coupled by hyperfine interaction, and it has zero rotational angular momentum. At a field of 1.9 mT, it has a binding energy of 5 MHz \times h, with respect to the F = 3, $m_F = 3$ two-atom asymptote (18). As one might expect, we find that optical transition rates as measured below are improved when using this effectively more deeply bound state as the initial state for two-photon transfer instead of state |s>. We shut off the trap and perform all subsequent experiments in free flight. This does not affect the particle density immediately but reduces it during the later detection procedure, which takes about 6 ms, to avoid collisions between atoms and weakly bound dimers and, hence, loss. We detect molecules in $|a\rangle$ via states $|s\rangle$ and $|d\rangle$ by first applying a magnetic field gradient for atom-molecule Stern-Gerlach separation, then reversing the magnetic field ramp, and finally dissociating them on the Feshbach resonance at 4.8 mT and imaging the resulting atoms (17).

Efficient two-photon transfer via the stimulated Raman adiabatic passage (STIRAP) technique (14, 19) relies on a suitable choice for the excited state |e>. In our case, this state must have singlet character so that it can be used as a transfer state to deeply bound levels of the $X^1\Sigma_g^+$ potential. In general, it must be well separated from other states, which otherwise could be off-resonantly excited. It should thus be situated far to the red of the excited $S_{1/2} + P_{1/2}$ potential asymptote to avoid the high density of excited molecular states near that asymptote. We have performed optical loss spectroscopy starting from state $|a\rangle$ in the wavelength range of 1120 to 1130 nm, \sim 2300 cm⁻¹ to the red of the cesium D_1 line. For this measurement, we recorded the number of remaining molecules in $|a\rangle$ as a function of excitation wavelength and found two progressions of lines, which we assign to the potential curves of the mixed $(A^1 \Sigma_u^+ - b^3 \Pi_u) 0_u^+$ excited states and to the $(1)^{3}\Sigma_{g}^{+}$ excited state, respectively. For the present experiments, we choose for $|e\rangle$ a level of the 0^+_u progression that is 8879.63(1) cm⁻¹ above the F = 3, $m_F = 3$ two-atom asymptote, corresponding to a transition wavelength of 1126.173(1) nm (Fig. 1A). We measure all wavelengths on a home-built wave meter. We identify this previously unknown level as the 225th one of the 0^+_{μ} system, with an uncertainty of two in the absolute numbering.

The ground state level $|g\rangle$ with vibrational quantum number v = 73 is well known from conventional molecular spectroscopy (20, 21). However, its binding energy, as well as the binding energy of all deeply bound vibrational levels, has only been known with an uncertainty of $\sim \pm 0.45$ cm⁻¹ before the present experiments (21). We search for $|g\rangle$ by simultaneously exciting the transition from $|a\rangle$ to $|e\rangle$ with laser L_1 and the one from $|e\rangle$ to $|g\rangle$ with laser L_2 . The two light fields create a molecule-molecule dark state. The molecules initially in $|a\rangle$ are lost unless the second laser L_2 is on two-photon resonance, provided that the Rabi frequency Ω_2 on the second transition is $\geq \Omega_1$, the Rabi frequency on the first transition. For coherence, stability, and reproducibility, we lock both lasers to independent narrow-band optical resonators, which we reference to an optical frequency comb (22). The comb is not calibrated but it allows precise differential frequency measurements and provides long-term stability needed

for systematic line searches (23). We find the resonance condition with vibrational level v = 73 at 1005.976(1) and 1005.982(1) nm, corresponding to rotational quantum numbers J = 0 and 2. Identification of *J* is possible because the rotational energy splitting is well known. Figure 2, A and B, shows typical molecular

dark resonances when we set L_2 on resonance and step the detuning Δ_1 of L_1 near 1126.173 nm. Figure 2C shows a dark resonance involving v =73, J = 2 using a different excited molecular state |e'>, which is excited with L_1 near 1123.104 nm. Figure 2, D to F, shows dark resonances involving the neighboring vibrational levels v =



Fig. 2. Dark resonances for vibrational levels v = 71, 72, and 73. Laser L_2 is held on resonance while the detuning Δ_1 of L_1 is scanned. We record the number of molecules in $|a\rangle$ while both lasers are pulsed on simultaneously. (**A** to **C**) Dark resonances involving v = 73 for excitation with L_1 near 1126 nm into J = 0 and 2 and for excitation with L_1 near 1123 nm into J = 2, respectively. (**D** to **F**) Neighboring levels v = 71 and 72 for excitation near 1123 nm. The solid line in (B) is the result of a three-level model calculation matched to the data giving $\Omega_1 = 2\pi \times 2$ kHz $\sqrt{I_1/(\text{mW/cm}^2)}$ and $\Omega_2 = 2\pi \times 11$ kHz $\sqrt{I_2/(\text{mW/cm}^2)}$ for a pulse time of 5 µs at intensities of $I_1 = 4 \times 10^5$ mW/cm² for L_1 and $I_2 = 2 \times 10^5$ mW/cm² for L_2 , assuming a laser linewidth of 2 kHz.



Fig. 3. STIRAP transfer from the weakly bound state $|a\rangle$ to the deeply bound state $|g\rangle = |v| = 73$, $J = 2\rangle$ and back to $|a\rangle$. (**A**) Number of molecules in state $|a\rangle$ as a function of STIRAP time τ for $\Delta_1 \approx 0 \approx \Delta_2$. The measured pulse overlap begins at 5 µs and ends at ~15 µs. The second pulse overlap starts at 25 µs and ends at ~33 µs. (**B**) Schematic of the timing for the Rabi frequencies Ω_i , i = 1, 2, during the double STIRAP sequence. Laser L_1 is left on after the first STIRAP sequence to clear out any remaining population in $|a\rangle$. (**C**) Double STIRAP efficiency as a function of the detuning Δ_2 of laser L_2 for $\Delta_1 \approx 0$. The solid line is a Gaussian fit with a FWHM of 811 kHz. The peak Rabi frequencies are $\Omega_1 \approx 2\pi \times 3$ MHz and $\Omega_2 \approx 2\pi \times 6$ MHz. Error bars refer to the 1 σ error in determining the particle number.

71 and 72. These $X^1\Sigma_{\sigma}^+$ levels were easily found based on previously acquired Cs2 spectra (21). We determine the binding energy of these levels, with respect to the atomic $F_1 = 3$, $F_2 = 3$ asymptote at zero magnetic field, to be 1060.9694(10), 1088.3101(10), and 1115.9148(10) cm^{-1} for v = 73, 72, and 71 with J = 0, respectively. The binding energy of the rovibrational ground state v = 0 is thus 3628.7053(14) cm⁻¹, which represents an improvement in precision of more than two orders of magnitude compared with the previous determination (21). Fitting the data for the dark resonances with a three-level model taking into account offresonant excitations and laser linewidths, we determine the molecular transition strengths as given by the normalized Rabi frequencies for the transitions $|a\rangle$ to $|e\rangle$ and $|e\rangle$ to $|v = 73, J = 2\rangle$ to be $\Omega_1 = 2\pi \times 2 \text{ kHz } \sqrt{I/(\text{mW/cm}^2)}$ and $\Omega_2 = 2\pi \times 11 \text{ kHz } \sqrt{I/(\text{mW/cm}^2)}$, respectively. A comparison with a typical atomic transition strength of $\Omega_a = 2\pi \times 5$ MHz $\sqrt{I/(mW/cm^2)}$ giving $|\Omega_1/\Omega_a|^2 < 10^{-6}$ reflects the minuteness of the wave function overlap.

We are now in a position to carry out coherent transfer using the STIRAP technique. For $|g\rangle$ we choose the vibrational level with v =73, J = 2. This level will allow us to reach the rovibrational ground state v = 0, J = 0 with a second STIRAP step in view of the selection rule $\Delta J = 0, \pm 2$. STIRAP uses a counterintuitive overlapping pulse sequence in which L_2 is pulsed on before L_1 . As is well known (19), STIRAP relies on the existence of a dark state of the form $|D\rangle = \alpha(t)|a\rangle + \beta(t)|g\rangle$ (here, $\alpha(t)$ and $\beta(t)$ are time-dependent amplitudes with $|\alpha(t)|^2 + |\beta(t)|^2 = 1$). With sufficient adiabaticity, the function $|\alpha(t)|^2$ decreases smoothly from 1 to 0, while the function $|\beta(t)|^2$ increases smoothly from 0 to 1. The initial state $|a\rangle$ is thus rotated via $|D\rangle$ into the final state $|g\rangle$. The criterion for adiabaticity is $\tau_p \Omega^2 \gg (2\pi)^2 \Gamma$, where τ_p is the pulse overlap time, $\Omega \approx \Omega_1 \approx \Omega_2$ is the peak Rabi frequency during the pulse, and

 $\Gamma \approx 2\pi \times 4$ MHz is the (spontaneous) decay rate from the upper state $|e\rangle$, as determined from our loss measurements. This criterion is quite stringent, in particular, in view of the low wave function overlap that enters into Ω . An upper (experimental) limit for τ_p is given by the relative laser coherence time for L_1 and L_2 . We choose τ_p to be ~10 µs. For detection, we apply the reverse STIRAP sequence after a waiting time $\tau_w \approx 10 \,\mu$ s to transfer the molecules back into $|a\rangle$. During this time we leave laser L_1 on to assure that all possible residual population in state $|a\rangle$ is removed.

We perform double STIRAP ~3 ms after the production of the Feshbach molecules and 1 ms after shutting off the trap. Figure 3A shows the molecular population in $|a\rangle$ as a function of the STIRAP time τ , and Fig. 3B shows the timing sequence for the double-transfer scheme. For recording the time evolution of the population, we interrupt the transfer process after time τ and measure the remaining population in $|a\rangle$. The molecules in $|a\rangle$ initially disappear during the first STIRAP sequence. They are now in level |v = 73, J = 2> of the singlet $X^1 \Sigma_g^+$ potential. Then a large fraction of them returns in the course of the reverse STIRAP sequence. For this particular measurement both lasers are on resonance. The peak Rabi frequencies are $\Omega_1 \approx$ $2\pi \times 3$ MHz and $\Omega_2 \approx 2\pi \times 6$ MHz. We typically obtain an overall efficiency >65% for the doubletransfer process, corresponding to single-pass efficiencies >80%, assuming equal efficiencies for both passes. Figure 3C shows the double-pass efficiency as a function of detuning Δ_2 of laser L_2 . Simulations for the three-level system show that the ~800-kHz full width at half maximum (FWHM) of the efficiency curve is compatible with a combination of laser power broadening and Fourier broadening. Our simulations also show that higher transfer efficiencies can be expected for an optimized STIRAP pulse sequence in which both peak Rabi frequencies are equal. Molecules not transferred by STIRAP are resonantly excited to



Fig. 4. (**A**) Ramsey-type experiment. The population in the initial state $|a\rangle$ oscillates as the hold time τ_h (during which both transfer lasers are off) is increased. The solid line is a sinusoidal fit to the data up to $\tau_h = 20 \ \mu$ s. Its frequency *f* is 115(2) kHz, in good agreement with the expected value of 113 kHz. The thin lines are borders to a region that is given by varying *f* by $\pm 6 \ \text{kHz}$, illustrating the estimated jitter in the two-photon detuning $|\Delta_2 - \Delta_1|$. (**B**) Comparison of the rate of expansion in the horizontal direction for the molecular sample without and with STIRAP transfer. The top curve (circles) shows the Thomas-Fermi radius *r* of the molecular sample as a function of expansion time without STIRAP. The linear fit gives a rate of expansion of *dr/dt* = 1.0(1) mm/s, corresponding to an energy of $k_B \times 14(4)$ nK. The bottom curve (squares) shows the expansion after double STIRAP with *dr/dt* = 0.7(1) mm/s, corresponding to $k_B \times 7(2)$ nK. Error bars indicate the 1σ error in the determination of the cloud radius.

le> and then lost from our three-level system by spontaneous emission into a multitude of ground state levels.

We demonstrate coherence of the transfer process in a Ramsey-type experiment (14), halting the transfer process by simultaneously shutting off both lasers 12 µs into the first STIRAP sequence when a balanced superposition of |a> and |g> has been created with $|\alpha(\tau)|^2 \approx 1/2 \approx |\beta(\tau)|^2$. After a hold time τ_h we resume the STIRAP transfer, with the roles of lasers L_1 and L_2 reversed. Thus, for $\tau_h = 0$ the population will simply be rotated back into the initial state. A three-level calculation shows that the population in the initial state $|a\rangle$ is expected to oscillate at the rate of the two-photon detuning $|\Delta_2 - \Delta_1|/(2\pi)$. Figure 4A shows the initial state population for $\Delta_1 \approx 0$ and $\Delta_2 \approx 2\pi \times 113$ kHz as a function of τ_h . The population oscillates at a frequency of $|\Delta_2 - \Delta_1|/(2\pi)$, however with marked increase in phase jitter on the time scale of 30 us. We attribute this apparent loss of phase coherence to a slow relative frequency drift of lasers L_1 and L_2 , leading to a slightly different twophoton detuning from one experimental run to the next. In Fig. 4A, we have added a region indicating a frequency jitter of ±6 kHz. This value is compatible with the present long-term stability of our lasers. The frequency drift does not affect an individual STIRAP process because the transfer efficiency is very robust against laser detuning, as shown in Fig. 3C.

We now show that the molecular sample is not heated during the transfer process and is indeed in the quantum gas regime. Specifically, we measure and compare the rate of expansion of the molecular sample in state $|a\rangle$ with and without the double-transfer process. In our regime, the energy scale for expansion is usually set by the mean field of the BEC, resulting in typical expansion energies for the atoms in the range from $k_{\rm B} \times 2$ nK to $k_{\rm B} \times 10$ nK (where $k_{\rm B}$ is Boltzmann's constant), depending on the strength of the atomic interaction (24). We find that the initial magnetic field ramping excites collective motion of the BEC in the form of a breathing mode as a result of a change in the mean field potential due to a change in atomic interaction strength (16). The breathing is transformed into expansion of the sample when the trap is shut off. We follow the expansion by monitoring the change of the Thomas-Fermi radius r of the sample. Figure 4B shows this radius along the horizontal direction as a function of expansion time with and without STIRAP. Without STIRAP, we obtain from a linear fit an expansion rate of dr/dt = 1.0(1) mm/s, corresponding to an energy of $k_{\rm B} \times 14(4)$ nK. With STIRAP, the rate is dr/dt = 0.7(1) mm/s, corresponding to an energy of $k_{\rm B} \times 7(2)$ nK. Both values are compatible with a separate measurement of the expansion of the atomic BEC for the same magnetic field ramp. Interestingly, the rate for the case with STIRAP is lower. We speculate that STIRAP with the

REPORTS

tightly focused laser beams L_1 and L_2 preferentially transfers molecules in the center of the sample and is hence responsible for some selection in velocity space.

It should now be possible to add a second STIRAP step for transfer into the rovibrational ground state v = 0, J = 0. A suitable two-photon transition at readily available laser wavelengths is via the 68th excited state level of the 0^+_{μ} potential near 1329 nm (up) and 991 nm (down) with comparatively good wave function overlap at the level of $|\Omega/\Omega_a|^2 \approx 10^{-4}$. We expect that searching for dark resonances will be straightforward, as now all two-photon transition energies are known with an uncertainty of 10^{-3} cm⁻¹. Molecules in v = 0, J = 0 cannot further decay into a lower state upon a two-body collision, and they are thus expected to allow the formation of an intrinsically stable molecular BEC. The high speed of our STIRAP transfer will allow us to perform in situ as well as time-of-flight imaging for direct characterization of the spatial and momentum distribution of the molecular ensemble.

With our technique, any low-lying vibrational state can be coherently populated in a controlled fashion with full control over the rotational quantum number, allowing, for instance, state-specific collisional studies and high-precision molecular spectroscopy with possible implications for fundamental physics (6, 7). Our procedure can be adapted to other species, in particular to heteronuclear alkali dimers such as RbCs (25) and KRb (15) for the creation of dipolar quantum gases (26). For heteronuclear alkali dimers, a single two-photon transfer step might suffice as

a result of favorable wave function overlap (27). We expect that the combination of our technique with Feshbach molecule production out of a Mottinsulator state in a three-dimensional lattice (28)will increase the initial Feshbach molecule production efficiency, avoiding collective excitations as a result of magnetic field ramping and inhibiting collisional loss, and will provide full control over all internal and external quantum degrees of freedom of the ground state molecules.

References and Notes

- 1. C. Chin et al., Phys. Rev. Lett. 94, 123201 (2005).
- 2. T. Kraemer et al., Nature 440, 315 (2006).
- 3. P. Staanum, S. D. Kraft, J. Lange, R. Wester, M. Weidemüller, Phys. Rev. Lett. 96, 023201 (2006).
- 4. N. Zahzam, T. Vogt, M. Mudrich, D. Comparat, P. Pillet, Phys. Rev. Lett. 96, 023202 (2006).
- 5. R. V. Krems, Int. Rev. Phys. Chem. 24, 99 (2005).
- 6. T. Zelevinsky, S. Kotochigova, J. Ye, Phys. Rev. Lett. 100, 043201 (2008).
- 7. D. DeMille et al., Phys. Rev. Lett. 100, 043202 (2008). M. Inguscio, W. Ketterle, C. Salomon, Eds., Ultracold Fermi
- Gases, Proceedings of the International School of Physics "Enrico Fermi," Course CLXIV (IOS Press, Amsterdam, 2008). D. DeMille, Phys. Rev. Lett. 88, 067901 (2002).
- 10.]. Doyle, B. Friedrich, R. V. Krems, F. Masnou-Seeuws,
- Eur. Phys. I. D 31, 149 (2004). T. Köhler, K. Góral, P. S. Julienne, Rev. Mod. Phys. 78, 11.
- 1311 (2006). 12. K. M. Jones, E. Tiesinga, P. D. Lett, P. S. Julienne, Rev.
- Mod. Phys. 78, 483 (2006). 13. D. Jaksch, V. Venturi, J. I. Cirac, C. J. Williams, P. Zoller,
- Phys. Rev. Lett. 89, 040402 (2002). 14. K. Winkler et al., Phys. Rev. Lett. 98, 043201 (2007).
- 15. S. Ospelkaus et al., preprint available at http://arxiv.org/ abs/0802 1093 (2008)
- T. Weber, J. Herbig, M. Mark, H.-C. Nägerl, R. Grimm, 16 Science 299, 232 (2003), published online 5 December 2002: 10.1126/science.1079699.

- 17.]. Herbig et al., Science 301, 1510 (2003), published online 21 August 2003; 10.1126/science.1088876.
- 18. M. Mark et al., Phys. Rev. A 76, 042514 (2007). K. Bergmann, H. Theuer, B. W. Shore, Rev. Mod. Phys.
- 70, 1003 (1998). 20. W. Weickenmeier et al., J. Chem. Phys. 82, 5354 (1985).
- 21. C. Amiot, O. Dulieu, J. Chem. Phys. 117, 5155 (2002).
- 22. Laser L_1 near 1126 nm and laser L_2 near 1006 nm are continuous-wave grating-stabilized tunable diode lasers with up to 26 mW and 5 mW of power at the sample position, respectively, both focused to a $1/e^2$ waist of ~25 μ m for sufficiently high Rabi frequencies. We estimate the laser linewidth for both lasers to be on the order of 1 kHz. The laser beams propagate horizontally at an angle of 80° with the long axis of the BEC with vertical linear polarization. Copropagation assures that the imparted photon recoil during STIRAP is minimal, corresponding to an energy of $k_{\rm B} \times 0.4$ nK, with Boltzmann's constant $k_{\rm B}$. The beam intensity is controlled by acousto-optical modulators, allowing pulse lengths down to 1 $\ensuremath{\mu s}$.
- 23. The wave-meter calibration is currently not sufficient to allow absolute numbering of the frequency comb teeth.
- 24. T. Kraemer et al., Appl. Phys. B 79, 1013 (2004).
- 25. J. M. Sage, S. Sainis, T. Bergeman, D. DeMille, Phys. Rev. Lett. 94, 203001 (2005).
- 26. K. Góral, L. Santos, M. Lewenstein, Phys. Rev. Lett. 88, 170406 (2002).
- W. C. Stwalley, *Eur. Phys. J. D* **31**, 221 (2004).
 T. Volz *et al.*, *Nat. Phys.* **2**, 692 (2006).
- 29. We thank the team of]. Hecker Denschlag, the LevT team in our group, and T. Bergeman for very helpful discussions and M. Prevedelli for technical assistance. We are indebted to R. Grimm for generous support and gratefully acknowledge funding by the Austrian Ministry of Science and Research (Bundesministerium für Wissenschaft und Forschung) and the Austrian Science Fund (Fonds zur Förderung der wissenschaftlichen Forschung) in the form of a START prize grant and by the European Science Foundation in the framework of the EuroQUAM collective research project QuDipMol.

1 May 2008; accepted 1 July 2008 Published online 10 July 2008; 10.1126/science.1159909 Include this information when citing this paper.

Observation of Atomic Diffusion at Twin-Modified Grain **Boundaries in Copper**

Kuan-Chia Chen,¹* Wen-Wei Wu,²* Chien-Neng Liao,¹† Lih-Juann Chen,¹ K. N. Tu³

Grain boundaries affect the migration of atoms and electrons in polycrystalline solids, thus influencing many of the mechanical and electrical properties. By introducing nanometer-scale twin defects into copper grains, we show that we can change the grain-boundary structure and atomic-diffusion behavior along the boundary. Using in situ ultrahigh-vacuum and high-resolution transmission electron microscopy, we observed electromigration-induced atomic diffusion in the twin-modified grain boundaries. The triple point where a twin boundary meets a grain boundary was found to slow down grain-boundary and surface electromigration by one order of magnitude. We propose that this occurs because of the incubation time of nucleation of a new step at the triple points. The long incubation time slows down the overall rate of atomic transport.

rain boundaries affect many physical properties of polycrystalline solids. For example, reduction of grain size is known to improve the mechanical strength of metals, governed by the Hall-Petch equation (1, 2). A large-angle tilt-type grain boundary can shortcircuit atomic diffusion, which has been the most serious reliability issue in Al interconnects in microelectronics technology. The atomic structure of a grain boundary is controlled by the misorientation between the two grains forming the grain boundary. Balluffi et al. have made bicrystals of Au thin films and varied systematically the tilt or twist angle in the bicrystals for studying the correlation between formation energy and atomic structure of grain boundaries (3, 4).

Generally speaking, the higher the grain-boundary misorientation angle, the higher the atomic diffusivity. Thus, by modifying the structure of a grain boundary, it should be possible to control the atomic diffusion along the grain boundary.

Lu et al. have synthesized a high density of nanotwins in pure Cu foils by pulsed electrodeposition (5). The average grain size in the Cu foils is ~400 nm, and the high-density twins have a peak at 15 nm in twin-lamella size distribution. The Cu foil shows a 10-fold improvement of the mechanical strength relative to a large-grained Cu, and the foil remains ductile but its electrical resistance did not significantly change. High mechanical strength and low electrical resistivity are desired properties for interconnecting wires in integrated circuits from the consideration of the resistive-capacitive delay, electromigration (EM), and stress migration (6-8). EM is enhanced

*These authors contributed equally to this work. †To whom correspondence should be addressed. E-mail: cnliao@mx.nthu.edu.tw

¹Department of Materials Science and Engineering, National Tsing Hua University, Hsinchu 30013, Taiwan, Republic of China. ²Department of Materials Science and Engineering, National Chiao Tung University, Hsinchu 30010, Taiwan, Republic of China. ³Department of Materials Science and Engineering, University of California, Los Angeles, CA 90095, USA.

Dark resonances for ground state transfer of molecular quantum gases

Manfred J. Mark¹, Johann G. Danzl¹, Elmar Haller¹, Mattias Gustavsson¹, Nadia Bouloufa², Olivier Dulieu², Houssam Salami³, Tom Bergeman³, Helmut Ritsch⁴, Russell Hart¹, Hanns-Christoph Nägerl¹

- ² Laboratoire Aimé Cotton, CNRS, Université Paris-Sud, Bât. 505, 91405 Orsay Cedex, France
- ³ Department of Physics and Astronomy, SUNY Stony Brook, NY 11794-3800, USA
- ⁴ Institut für Theoretische Physik und Zentrum für Quantenphysik, Universität Innsbruck, Technikerstraße 25, A–6020 Innsbruck, Austria

Received: date / Revised version: date

Abstract One possible way to produce ultracold, highphase-space-density quantum gases of molecules in the rovibronic ground state is given by molecule association from quantum-degenerate atomic gases on a Feshbach resonance and subsequent coherent optical multi-photon transfer into the rovibronic ground state. In ultracold samples of Cs_2 molecules, we observe two-photon dark resonances that connect the intermediate rovibrational level $|v=73, J=2\rangle$ with the rovibrational ground state |v=0, J=0> of the singlet $X^1 \Sigma_g^+$ ground state potential. For precise dark resonance spectroscopy we exploit the fact that it is possible to efficiently populate the level |v=73, J=2 > by two-photon transfer from the dissociation threshold with the stimulated Raman adiabatic passage (STIRAP) technique. We find that at least one of the two-photon resonances is sufficiently strong to allow future implementation of coherent STIRAP transfer of a molecular quantum gas to the rovibrational ground state |v=0, J=0>.

1 Introduction

Laser cooling of atoms and the production of quantum degenerate atomic Bose and Fermi gases have revolutionized the field of atomic physics [1]. For molecular systems, ultralow temperatures and high phase space densities are much more difficult to achieve. Laser cooling of molecules has not yet been demonstrated, and with alternative cooling and slowing techniques such as buffer gas cooling and Zeeman slowing high phase space densities are yet out of reach [2–4]. In photoassociation experiments from magneto-optical traps, [5–9], cold samples of deeply bound molecules in the lowest vibrational levels have been created. Yet, the phase space densities are far away from the quantum degenerate regime. In the limit of extremely weak binding, molecular Bose-Einstein condensation could be achieved [10] by using the trick of first cooling an atomic Fermi gas to high phase space densities and subsequently associating pairs of atoms to molecules. For molecules composed of Fermions, collisional stability of the highly excited molecules is assured as a result of a Pauli blocking effect. Here, we are interested in ultracold and dense molecular systems in specific deeply bound rovibrational levels. Such samples are of high interest for fundamental studies in physics and chemistry, ranging from ultracold chemistry [11] and few-body collisional physics [12,13] to high resolution spectroscopy [14,15], to applications in quantum processing [16], and to the formation of dipolar quantum gases and dipolar Bose-Einstein condensates [17,18]. For these experiments full control over the molecular wave function is desired. In addition, high densities are required for molecular quantum gas studies. Only in the rovibronic ground state, i.e. the lowest energy level of the electronic ground state, is collisional stability assured.

For the production of molecular quantum gases in the absolute ground state, we follow a scheme in which the technique of stimulated two-photon transfer is repeatedly applied to molecules associated on a Feshbach resonance from a high-density sample of ultracold atoms such as a Bose-Einstein condensate (BEC). The initially very loosely bound molecules are to be transferred in a few successive steps to the rovibrational ground state, acquiring more and more binding energy. The scheme has several advantages. It is fully coherent, not relying on spontaneous processes, allowing high state selectivity, and it involves only a comparatively small number of intermediate levels. The scheme is expected to allow the removal of a ground state binding energy of typically 0.5 eV for an alkali dimer without appreciably heat-

¹ Institut für Experimentalphysik und Zentrum für Quantenphysik, Universität Innsbruck, Technikerstraße 25, A–6020 Innsbruck, Austria

Correspondence to: johann.danzl@uibk.ac.at

ing the molecular sample. It essentially preserves phase space density and coherence of the particle wave function, allowing the molecular sample to inherit the high initial phase space density from the atomic sample. Ideally, the scheme will ultimately result in the formation of a molecular BEC. A major challenge is given by the low radial wave function overlap between successive molecular levels, potentially leading to prohibitively low transition rates for the two-photon transitions that could only be compensated by the use of further (smaller) transfer steps.

In a crucial experiment, Winkler et al. [19] demonstrated that coherent two-photon transfer by means of the stimulated Raman adiabatic passage (STIRAP) technique [20] can efficiently be implemented with quantum gases of weakly bound Feshbach molecules. In this work, the transferred molecules, in this case Rb_2 , were still weakly bound with a binding energy of much less than 10^{-4} of the binding energy of the rovibrational ground state. In particular, wave function overlap of the final level with the rovibrational ground state is negligible. Nevertheless, an important result of this experiment was the demonstration that, even with excitation near the excited S+P asymptote, parasitic excitation of unwanted molecular transitions by the STIRAP laser beams could largely be avoided. Recently, Danzl et al. [21] showed efficient coherent STIRAP transfer into deeply bound rovibrational levels in the quantum gas regime. More specifically, transfer into the rovibrational level |v=73, J=2>of the singlet $X^1 \varSigma_g^+$ molecular potential of the Cs dimer was demonstrated. This level is bound by 1061 wavenumbers, more than one-fourth of the binding energy of the rovibrational ground state. Here, as usual, v and J denote the vibrational and rotational quantum numbers, respectively. This intermediate level was chosen as to give a balanced distribution for the wave function overlap in a four-photon transfer scheme to the ground state, i.e. to assure that all four dipole transition moments are of comparable magnitude. This level could thus serve as a transfer state towards the rovibrational ground state |v = 0, J = 0 >, allowing coherent ground state transfer with two two-photon transitions. Also recently, Ni et al. [22] could demonstrate transfer all the way into the rovibrational ground state |v=0, J=0 > of the singlet $X^1 \Sigma^+$ molecular potential in a quantum gas of KRb molecules. The transfer could be achieved in a single step as a result of the favorable run of the excited state potentials in the case of heteronuclear alkali dimers [23]. Also, the lowest rovibrational level of the Rb₂ triplet $a^3 \Sigma_{\mu}^+$ potential could recently be populated in the quantum gas regime using the STIRAP technique [24].

Here, in an ultracold and dense sample of Cs molecules, we present two-photon dark resonances connecting the rovibrational level |v = 73, J = 2 > of the Cs dimer singlet $X^1 \Sigma_g^+$ molecular potential with the rovibrational ground state |v = 0, J = 0 >. Starting from |v = 73, J = 2 >, we first perform molecular loss spec-



Fig. 1 Molecular level scheme for Cs₂. Molecules in a weakly bound Feshbach level $|1\rangle = |v \approx 155 \rangle$ (not resolved near the 6S+6S asymptote) are transferred to the rovibrational level $|3\rangle = |v = 73, J = 2\rangle$ of the singlet $X^1 \Sigma_g^+$ ground state potential with a binding energy of 1061 cm^{-1} by a twophoton STIRAP process [21] involving lasers L_1 and L_2 near 1126 nm and 1006 nm. The following two-photon transition from |3> to |5>=|v=0, J=0> and also to |v=0, J=2> is then probed by lasers L_3 and L_4 near 1350 nm and 1000 nm, respectively. Level $|2\rangle$ is the 225th level of the electronically excited coupled $(A^1 \Sigma_u^+ - b^3 \Pi_{0u}) 0_u^+$ potentials. Here, we probe suitable candidate levels for $|4\rangle$, connecting $|3\rangle$ to $|5\rangle$. These candidate levels also belong to the 0^+_{μ} coupled state system and include levels with coupled channel vibrational numbers v' = 57 to 68. The position of the vertical arrows is not meant to reflect the internuclear distance at which the transition takes place.

troscopy by laser excitation in the wavelength range from 1329 nm to 1365 nm to search for and identify suitable excited state levels of the mixed $(A^1 \Sigma_u^+ - b^3 \Pi_{0u}) 0_u^+$ excited molecular potentials. These levels are 9893 to 10091 wavenumbers above the rovibronic ground state, corresponding to a wavelength range from 1011 nm to 991 nm for the transition to the rovibronic ground state. We then perform dark state spectroscopy by simultaneous laser irradiation near 1350 nm and 1000 nm. We find several dark resonances, from which we derive normalized transition strengths and find that at least one of the two-photon transitions is favorable for ground state transfer.

2 Molecular energy levels and laser transitions

Fig.1 shows the energy of the relevant Cs₂ molecular states and the optical transitions for our transfer scheme. State $|1\rangle$ is the initial weakly bound Feshbach state that we populate out of an atomic BEC of Cs atoms via Feshbach association [25]. For the transfer from $|1\rangle$ to the ro-vibrational ground state $|5\rangle = |v = 0, J = 0\rangle$, three intermediate levels $|2\rangle$, $|3\rangle$, and $|4\rangle$ are needed. All five molecular levels are coupled by two two-photon transitions in a distorted M-shaped configuration as shown in



Fig. 2 5-level distorted M-scheme. The one-photondetunings and Rabi frequencies of L_i are Δ_i and Ω_i , i = 1, 2, 3, 4. For STIRAP to |v = 73, J = 2 > the detunings for L_1 and L_2 are $\Delta_1 \approx 0 \approx \Delta_2$.

Fig.2. Levels $|2\rangle$ and $|4\rangle$ belong to the excited mixed $(A^1 \Sigma_u^+ - b^3 \Pi_{0u}) 0_u^+$ potentials. We have identified level $|2\rangle$ as the 225th one of the coupled 0_u^+ system, with an uncertainty of 2 in the absolute numbering, and $|3\rangle$ is the level with v = 73 and J = 2 of the $X^1 \Sigma_q^+$ ground state potential [21]. A two-photon laser transition with laser L_1 at 1126 nm and laser L_2 at 1006 nm couples $|1\rangle$ to $|3\rangle$ via $|2\rangle$. There are now several possibilities for coupling $|3\rangle$ to $|5\rangle$, differing in the choice of the excited state |4>. The aim of this work is to identify a suitable state $|4\rangle$ from the $(A^1 \Sigma_u^+ - b^3 \Pi_{0u}) \ 0_u^+$ potentials with sufficient wave function overlap with both $|3\rangle$ and $|5\rangle$. We search for state $|4\rangle$ in the energy range of 9893 to 10091 wavenumbers above the rovibrational ground state |5>. Molecular structure calculations as outlined in Sec. 4 show that in this range there are candidate states for $|4\rangle$ that have dipole transition matrix elements with both $|3\rangle$ and $|5\rangle$ of comparable magnitude, allowing optimum STIRAP performance. The wavelengths for the lasers L_3 and L_4 driving the associated two-photon transition are near 1350 nm and 1000 nm, respectively. We derive all laser light for driving the molecular transitions from highly stable, widely tunable diode laser systems with kHz linewidths. For short term stability, the lasers are all locked to narrow-band optical resonators. For long term stability, the optical resonators are referenced to an infrared, fiber-laser-based frequency comb, covering the wavelength range from about 980 nm to about 2000 nm.

3 Preparation of a molecular quantum gas in v = 73, J = 2

Our sample preparation procedure follows Ref. [21]. In summary, we first produce a cigar-shaped BEC of typically 1.5×10^5 cesium atoms in the lowest hyperfine sublevel F=3, $m_F=3$ in a crossed optical dipole trap. As usual, F is the atomic angular momentum quantum number, and m_F its projection. The trapping light at 1064.5 nm is derived from a single-frequency, highlystable Nd:YAG laser. Using a d-wave Feshbach resonance at 4.8 mT [26] we then produce a quantum gas of weakly bound Feshbach molecules out of the BEC [25]. For this, we first ramp the magnetic field from the BEC production value of 2.0 mT to 4.9 mT, slightly above the Feshbach resonance. The molecules are produced on a downward sweep at a typical sweep rate of 0.025 mT/ms. The resulting ultracold sample contains up to 11000 molecules, immersed in the bath of the remaining BEC atoms. For the present experiments we shut off the trap and perform all subsequent measurements in free flight. This reduces the particle density, in particular during the later detection stage of the experiment, and hence reduces atom-molecule collisional loss, thus increasing the molecular signal. Following two avoided state crossings while further sweeping the magnetic field to lower values, we transfer the molecules via a weakly bound, open channel s-wave molecular state into the still weakly bound, closed channel s-wave molecular state $|1\rangle$ by magnetic field ramping [21]. This is the starting state for the subsequent optical transfer. As with all other weakly bound Feshbach states, it belongs to both the $X^1 \varSigma_g^+$ ground state potential and the lowest triplet $a^3 \Sigma_n^+$ potential and is hence of mixed character. It has zero rotational angular momentum. At a field of 1.9 mT, it has a binding energy of 5 MHz×h, where his Planck's constant, with respect to the $F = 3, m_F = 3$ two-atom asymptote [26]. We detect molecules in $|1\rangle$ by reverse magnetic field ramping, leading to dissociation on the Feshbach resonance at 4.8 mT, and by subsequent imaging of the resulting atoms [25].

We transfer the molecules from $|1\rangle$ to the rovibrational level $|3\rangle = |v = 73, J = 2\rangle$ with the STIRAP technique [21]. For this, about 3 ms after molecule production, with the magnetic field ramping completed, laser L_2 at 1006 nm is pulsed on first and then laser L_1 at 1126 nm. Both lasers are on resonance within a few kHz. The pulse overlap time is about 10 μ s. With peak Rabi frequencies of $\Omega_1 \approx 2\pi \times 3$ MHz and $\Omega_2 \approx 2\pi \times 6$ MHz we transfer about 80 % of the molecules to $|3\rangle$. We find that the molecular sample is not heated as a result of the STIRAP transfer. A residual kinetic energy on the order of $k_B \times 10$ nK comes from the expansion energy of the initial atomic sample. Our current procedure allows us to produce a sample of up to 8000 molecules in state $|3\rangle$ every 12 s. For the loss spectroscopy as detailed below, we irradiate the molecules in $|3\rangle$ with light near 1350 nm for a certain waiting time. We then measure the fraction of molecules that have remained in |3>. For this, we transfer the remaining molecules back to $|1\rangle$ using the reverse STIRAP process and determine the number of molecules in $|1\rangle$. Without irradiation with light near $1350~\mathrm{nm}$ we transfer more than 65% of the molecules from |1> to |3> and back to |1>[21].



Fig. 3 Loss resonances for excitation near 1351 nm from $|3\rangle = |v = 73, J = 2 \rangle$ of the $X^1 \Sigma_g^+$ ground state potential. (A) Loss of molecules in $|3\rangle$ as a function of laser detuning Δ_3 near 1351 nm after a waiting time of 20 μ s. The solid line represents a model calculation matched to the data yielding an excited state natural linewidth of $2\pi \times 2$ MHz. (B) Time dependence of molecular loss on resonance at 1351 nm for two different laser intensities. (1) $270 \pm 80 \text{ mW/cm}^2$, (2) 570 $\pm 80 \text{ mW/cm}^2$. The fitted exponential decay gives the decay constants $\tau = 26 \pm 4 \ \mu$ s for 270 mW/cm² and $\tau = 14 \pm 2 \ \mu$ s for 570 mW/cm².

4 Loss spectroscopy

Prior to the present experiments, the energies of the levels with predominant $A^1 \Sigma_u^+$ character in the region of interest were established to about $\pm 0.06 \text{ cm}^{-1}$ by fits [27] to data obtained by Fourier transform spectroscopy (FTS) at Laboratoire Aimé Cotton (LAC) using transitions to the $X^1 \Sigma_q^+$ state. However, the predominantly $b^3 \Pi_{0u}$ levels were only known to about $\pm 2 \text{ cm}^{-1}$ because this region was above that for which data was obtained from $2^3 \Delta_{1g} \to b^3 \Pi_{0u}$ emission [28], but lower than the regime where $b^3 \Pi_{0u}$ levels acquire sufficient singlet character (by spin-orbit mixing) to be observed in the FTS work. Paradoxically, the predominantly $b^3 \Pi_{0u}$ levels are of special interest here because they happen to have significant singlet character over regions of the internuclear distance that are most important for transitions of interest in this work.

The coupled channel calculations used to characterize the level structure of the strongly interacting $A^1 \Sigma_u^+$ and $b^3 \Pi_{0u}$ states employed methods developed from previous work on A and b states of K₂ [29,30], RbCs [31], Na₂ [32], and Rb₂ [33]. The DVR approach [34] was used to calculate eigenvalues primarily for two coupled channels, although some information on $b^3 \Pi_{1u}$ was found in the FTS data from LAC. Similar computational approaches, differing in the detailed numerical methods, have been applied recently also to the A and b states of NaRb [35].

Because of the initial $\pm 2 \text{ cm}^{-1}$ uncertainty in the positions of $b^3 \Pi_{0u}$ levels of interest, we decided to perform a systematic, broad-range search around expected transition energies in the wavelength range from 1329 nm to 1365 nm. For this, we perform double STIRAP from $|1\rangle$ to $|3\rangle$ and back with a waiting time of typically $\tau = 1$ ms. During the waiting time, we irradiate the sample with laser L_3 at an estimated intensity of $5 \cdot 10^4$ mW/cm². Laser L_3 is a diode laser with grating feedback. On the timescale of our experiment, the resonator of the laser is sufficiently stable, allowing systematic tuning of the laser without locking the laser to its external resonator. We step the laser frequency in units of typically 20 MHz by tuning the piezo element on the grating. We monitor the laser wavelength with a home-built wavemeter at approximately 300 MHz accuracy. For the initial broad range line search we increased the repetition rate of the experiment by stopping evaporative cooling slightly before condensation sets in. While stepping the laser, taking data points essentially at the cycle rate corresponding to the sample production time, we look for a dip in the molecule number. Once such a dip is found, typically consisting of a few data points, we perform a more precise scan by locking the laser to the external, highly-stable resonator and then the external resonator to the infrared frequency comb. This allows us to detune the laser with kHz precision. Fig.3 (A) shows a typical loss resonance near 1351 nm. We reduce the laser intensity such that on resonance at most 80% of the molecules are lost within 20 μ s. From such measurements the transition strength as given by the normalized Rabi frequency and the natural linewidth of the excited state can be deduced. The typical width of the excited state molecular levels that we have identified is $2\pi \times 2$ MHz, in agreement with typical expected lifetimes. Fig.3 (B) shows a measurement of the time dependence of the molecular loss. Here, we step the waiting time τ from 0 to 50 μ s, while the laser is kept on resonance. In total, we have found 7 excited levels belonging to the $(A^1 \Sigma_u^+ - b^3 \Pi_{0u}) \ 0_u^+$ coupled state system. They are listed in Table 1 along with the dominant overall character (either $A^1 \Sigma_u^+$ state or $b^3 \Pi_{0u}$ state) of the vibrational wave function as determined from the coupled state calculations. Within the wavelength range from 1329 nm to 1365 nm, theory predicts the existence of 5 more states of the 0^+_u coupled state system, whose energies are also displayed in Table 1. For most of them, the wave function overlap is not expected to be favorable for STIRAP transfer to X $^1\varSigma_g^+ ~|v\!=\!0>.$ However, an improved model of the energy level structure, based on all the data except one FTS point with a large residual, fits the observed transitions to a rms residual error of 0.02 cm^{-1} , indicating that additional resonances can be found with searches over very limited ranges of laser frequency.

5 Dark resonances with |v=0, J=0 > and |v=0, J=2 >

In our recent work [21] we could greatly improve the value for the binding energy of the rovibrational ground state $|5\rangle = |v = 0, J = 0\rangle$ by determining the binding energy of $|v = 73\rangle$ and using well-known data from conventional molecular spectroscopy [36,37]. Our measurement was limited by the calibration of our wavemeter,

not allowing us to determine the number of the teeth of the frequency comb, and by the precision of the spectroscopy data. Searching for $|5\rangle$ in dark state spectroscopy is now a straightforward task as only a range of about 0.002 wavenumbers needs to be scanned. We do this by exciting the transitions from $|3\rangle$ to $|4\rangle$ with laser L_3 and from $|4\rangle$ to $|5\rangle$ with laser L_4 simultaneously. The intensity for L_4 is typically $5 \cdot 10^4$ mW/cm^2 . As is well known, the two light fields create a molecule-molecule dark state. The molecules initially in $|3\rangle$ are lost unless laser L_4 is on two-photon resonance, provided that the Rabi frequency Ω_4 on the fourth transition is equal to or greater than Ω_3 , the Rabi frequency on the third transition. We look for the resonance condition with the rovibrational ground state |v=0, J=0>for some of the excited levels that we found above. Table 1 lists the observed transition wavelengths. We check that we can identify the level with rotational quantum number J = 2 as the rotational energy splitting is well known. Fig.4 shows typical molecular dark resonances when we set L_4 on resonance and step the detuning Δ_3 of L_3 near 1350 nm. From a three-level model matched to the data for the dark resonances, taking into account offresonant excitations and laser line widths, we determine the molecular transition strengths as given by the normalized Rabi frequencies. One of the two-photon transitions appears to be a particularly good candidate for STIRAP ground state transfer. It involves the excited state level $|4\rangle$ with vibrational number v' = 61 of the $(A^1 \Sigma_u^+ - b^3 \Pi_{0u}) 0_u^+$ coupled state system. For the transition from $|3\rangle$ to $|4\rangle$ and from $|4\rangle$ to $|5\rangle$ the normalized Rabi frequencies are $\Omega_3 = 2\pi \times 6 \text{ kHz} \sqrt{I/(\text{mW/cm}^2)}$ and $\varOmega_4\!=\!2\pi\!\times\!5~\mathrm{kHz}~\sqrt{I/(\mathrm{mW/cm}^2)},$ respectively. These values carry an estimated error of 50% as the laser beam parameters for L_3 and L_4 are not well determined. A comparison with a typical atomic transition strength of $\Omega_a = 2\pi \times 5 \text{ MHz} \sqrt{I/(\text{mW/cm}^2)}$ giving $|\Omega_3/\Omega_a|^2 \approx 10^{-6}$ and $|\Omega_4/\Omega_a|^2 \approx 10^{-6}$ reflects the minuteness of the wave function overlap. Nevertheless, their value is sufficient for STIRAP as seen in our recent work [21]. Also, they are of similar magnitude. This facilitates STIRAP, for which the peak Rabi frequencies should be approximately equal for optimum performance.

6 Conclusion

We observe several two-photon dark resonances that connect the intermediate rovibrational level |v=73, J=2 >of the $X^1 \Sigma_g^+$ ground state potential with the rovibrational ground state level |v=0, J=0 >. At least one of the two-photon transitions is sufficiently strong for implementing STIRAP to |v=0, J=0 > in the quantum gas regime, paving the way for the realization of a BEC of ground state molecules. STIRAP can in principle be

Fig. 4 Dark resonances involving $X^1 \Sigma_g^+$ state levels |v = 73, J=2 > and |v=0 > for two different intermediate levels. (A and B) Dark resonances with $X^1 \Sigma_g^+ |v=0, J=0 >$ and |v=0, J=2 > involving the 0_u^+ excited state level |v'=63, J=1 > at an excitation wavelength near 1345 nm. (C and D) Dark resonances with $X^1 \Sigma_g^+ |v=0, J=0 >$ and |v=0, J=2 > involving the excited state level |v'=61, J=1 > at an excitation wavelength near 1351 nm. The solid line in (D) is the result of a model calculation, solving the three-level master equation including laser bandwidth and loss, matched to the data giving $\Omega_3 = 2\pi \times 6$ kHz $\sqrt{I/(\text{mW/cm}^2)}$ and $\Omega_4 = 2\pi \times 4$ kHz $\sqrt{I/(\text{mW/cm}^2)}$ for $X^1 \Sigma_g^+ |v=0, J=2 >$. The corresponding calculation for $X^1 \Sigma_g^+ |v=0, J=0 >$ yields $2\pi \times 5$ kHz $\sqrt{I/(\text{mW/cm}^2)}$.

implemented in two ways, either in the form of two sequential two-photon STIRAP steps, or in the form of four-photon STIRAP [38,39]. An attractive strategy for the production of a BEC of ground state molecules relies on the addition of an optical lattice. Starting from an atomic BEC, pairs of atoms at individual lattice sites are produced in a superfluid-to-Mott-insulator transition [40]. These pairs can then be very efficiently associated on a Feshbach resonance and subsequently transferred to the rovibronic ground state with STIRAP. The lattice has the advantage of shielding the molecules against inelastic collisions during the association process and subsequent state transfer. As proposed by Jaksch et al. [41], dynamical melting of the lattice should ideally result in the formation of a BEC of molecules in the rovibronic ground state in a Mott-insulator-to-superfluid-type transition.

7 Acknowledgements

We are indebted to R. Grimm for generous support and we thank E. Tiemann for valuable discussions and C. Amiot for providing the FTS data of LAC on Cs_2 . We gratefully acknowledge funding by the Austrian Ministry of Science and Research (BMWF) and the Austrian Science Fund (FWF) in form of a START prize grant and by the European Science Foundation (ESF) **Table 1** Levels of the excited 0^+_u coupled state system in the region 9893 cm⁻¹ to 10091 cm⁻¹ above $X^1 \Sigma_g^+ | v = 0, J = 0 >$. The first column gives the coupled channel vibrational numbers of the individual levels. Levels marked with * have not been searched for and the level energies given are those determined from the coupled channels calculations. The column labeled 'C' gives the predominant contribution to the overall vibrational wave function, which is either predominantly $A^1 \Sigma_u^+$ or predominantly $b^3 \Pi_{0u}$, indicated by A and b, respectively. The number in brackets gives the order within the two progressions of levels with either predominantly $A^1 \Sigma_{\mu}^+$ or predominantly $b^3 \Pi_{0u}$ character. Both the $|J=1\rangle$ and the |J=3> rotational levels were identified for all observed excited state levels. The wavemeter accuracy gives a typical uncertainty in wavelength of $\pm 0.002 \,\mathrm{nm}$, which translates into $\pm 0.011 \,\mathrm{cm}^{-1}$ uncertainty in the value for the energy above |v=0, J=0>. The energy relative to $X^1 \Sigma_g^+$ |v=0, J=0>of experimentally determined levels is based on the measured excitation wavelength from $X^1 \Sigma_q^+ | v = 73, J = 2 >$ and the $X^{1}\Sigma_{q}^{+} | v = 73 >$ level energy from Ref. [37], which introduces an additional uncertainty of 0.001 cm^{-1} . Deexcitation wavelengths are obtained from dark resonance spectroscopy involving the respective intermediate excited state level and the rovibronic ground state $X^1 \Sigma_q^+ | v = 0, J = 0 >$. n. m.: not measured

v'	С	J	Excitation wavelength from $X^{1}\Sigma_{g}^{+} v =$ 73, J = 2 > [nm]	Energy above $X^{1}\Sigma_{g}^{+} v =$ 0, J = 0 > $[\text{cm}^{-1}]$	De- excitation wavelength to $X^1 \Sigma_g^+$ v = 0, J = 0 > [nm]
57	A (7)	1	1365.148	9893.002	n. m.
57	A (7)	3	1365.131	9893.094	n. m.
*58	b(50)	0	1362.893	9905.126	n. m.
*59	A (8)	0	1357.748	9932.927	n. m.
60	b(51)	1	1357.091	9936.497	n. m.
60	b(51)	3	1357.071	9936.606	n. m.
61	b(52)	1	1351.367	9967.707	1003.240
61	b(52)	3	1351.347	9967.816	n. m.
*62	A (9)	0	1350.388	9973.068	n. m.
63	b(53)	1	1345.725	9998.729	1000.128
63	b(53)	3	1345.705	9998.839	n. m.
*64	A(10)	0	1343.082	10013.351	n. m.
65	b(54)	1	1340.162	10029.576	997.052
65	b(54)	3	1340.143	10029.682	n. m.
66	A (11)	1	1335.833	10053.759	994.653
66	A (11)	3	1335.816	10053.853	n. m.
*67	b (55)	0	1334.675	10060.249	n. m.
68	b(56)	1	1329.257	10090.794	991.003
68	b (56)	3	1329.238	10090.902	n. m.

in the framework of the EuroQUAM collective research project QuDipMol. R.H. acknowledges support by the European Union in form of a Marie-Curie International Incoming Fellowship (IIF). The work at Stony Brook was supported by the US NSF, under grant PHY0652459.

M.J. Mark et al.

References

- K. Southwell (ed.), Ultracold matter, Nature (Insight) 416 205246 (2002).
- J. Doyle, B. Friedrich, R.V. Krems, and F. Masnou-Seeuws, Eur. Phys. J. D 31, 149 (2004).
- 3. R.V. Krems, Phys. Chem. Chem. Phys. 10, 4079 (2008).
- S.Y.T. van de Meerakker, H.L. Bethlem, and G. Meijer, Nature Phys. 4, 595 (2008).
- For a review on photoassociation, see: K. M. Jones, E. Tiesinga, P.D. Lett, and P. S. Julienne, Rev. Mod. Phys. 78, 483 (2006).
- A.N. Nikolov, J.R. Ensher, E.E. Eyler, H. Wang, W.C. Stwalley, and P.L. Gould, Phys. Rev. Lett. 84, 246 (2000).
- J.M. Sage, S. Sainis, T. Bergeman, and D. DeMille, Phys. Rev. Lett. 94, 203001 (2005).
- M. Viteau, A. Chotia, M. Allegrini, N. Bouloufa, O. Dulieu, D. Comparat, and P. Pillet, Science **321**, 232 (2008).
- J. Deiglmayr, A. Grochola, M. Repp, K. Mörtlbauer, C. Glück, J. Lange, O. Dulieu, R. Wester, and M. Weidemüller, Phys. Rev. Lett. **101**, 133004 (2008).
- For an overview, see: Ultracold Fermi Gases, Proceedings of the International School of Physics Enrico Fermi, Course CLXIV, edited by M. Inguscio, W. Ketterle, and C. Salomon (IOS Press, Amsterdam, 2008).
- 11. R. V. Krems, Int. Rev. Phys. Chem. 24, 99 (2005).
- P. Staanum, S. D. Kraft, J. Lange, R. Wester, and M. Weidemüller, Phys. Rev. Lett. 96, 023201 (2006).
- N. Zahzam, T. Vogt, M. Mudrich, D. Comparat, and P. Pillet, Phys. Rev. Lett. 96, 023202 (2006).
- T. Zelevinsky, S. Kotochigova, and J. Ye, Phys. Rev. Lett. 100, 043201 (2008).
- D. DeMille, S. Sainis, J. Sage, T. Bergeman, S. Kotochigova, and E. Tiesinga, Phys. Rev. Lett. **100**, 043202 (2008).
- 16. D. DeMille, Phys. Rev. Lett. 88, 067901 (2002).
- K. Góral, L. Santos, and M. Lewenstein, Phys. Rev. Lett. 88, 170406 (2002).
- M. Baranov, L. Dobrek, K. Góral, L. Santos, and M. Lewenstein, Phys. Scr. **T102**, 74 (2002).
- K. Winkler, F. Lang, G. Thalhammer, P. v. d. Straten, R. Grimm, and J. Hecker Denschlag, Phys. Rev. Lett. 98, 043201 (2007).
- K. Bergmann, H. Theuer, and B. W. Shore, Rev. Mod. Phys. 70, 1003 (1998).
- J. G. Danzl, E. Haller, M. Gustavsson, M.J. Mark, R. Hart, N. Bouloufa, O. Dulieu, H. Ritsch, and H.-C. Nägerl, Science **321**, 1062 (2008), published online 10 July 2008; 10.1126/science.1159909.
- K.-K. Ni, S. Ospelkaus, M. H. G. de Miranda, A. Peer, B. Neyenhuis, J.J. Zirbel, S. Kotochigova, P.S. Julienne, D.S. Jin, and J. Ye, published online September 18, 2008, Science DOI: 10.1126/science.1163861.
- 23. W.C. Stwalley, Eur. Phys. J. D **31**, 221 (2004).
- 24. F. Lang, K. Winkler, C. Strauss, R. Grimm, and J. Hecker Denschlag, Phys. Rev. Lett. **101**, 133005 (2008).
- 25. J. Herbig, T. Kraemer, M. Mark, T. Weber, C. Chin, H.-C. Nägerl, and R. Grimm, Science **301**, 1510 (2003); published online 21 August 2003; 10.1126/science.1088876.
- M. Mark, F. Ferlaino, S. Knoop, J.G. Danzl, T. Kraemer, C. Chin, H.-C. Nägerl, and R. Grimm, Phys. Rev. A 76, 042514 (2007).

Dark resonances for ground state transfer of molecular quantum gases

- 27. H. Salami, T. Bergeman, O. Dulieu, D. Li, F. Xie, and L. Li, manuscript in preparation (2008).
- 28. F. Xie, D. Li, L. Tyree, L. Li, V. Sovkov, V. S. Ivanov, S. Magnier, and A. M. Lyyra, J. Chem. Phys. **128**, 204313 (2008).
- 29. C. Lisdat, O. Dulieu, H. Knöckel and E. Tiemann, Eur. Phys. J. D 17, 319 (2001).
- 30. M. R. Manaa, A. J. Ross, F. Martin, P. Crozet, A. M. Lyyra, L. Li, C. Amiot and T. Bergeman, J. Chem. Phys. 117, 11208 (2002).
- 31. T. Bergeman, C. E. Fellows, R. F. Gutterres and C. Amiot, Phys. Rev. A, 67, 050501 (2003).
- 32. P. Qi, J. Bai, E. Ahmed, A. M. Lyyra, S. Kotochigova, A. J. Ross, C. Effantin, P. Zalicki, J. Vigué, G. Chawla, R. W. Field, T.-J. Whang, W. C. Stwalley, H. Knöckel, E. Tiemann, J. Shang, L. Li, and T. Bergeman, J. Chem. Phys. **127**, 044301 (2007).
- 33. H. Salami, T. Bergeman, B. Beser, J. Bai, E. Ahmed, S. Kotochigova, A. M. Lyyra, J. Huennekens, C. Lisdat, O. Dulieu, P. Crozet and A. J. Ross, manuscript in preparation (2008).
- 34. D. Colbert and W. H. Miller, J. Chem. Phys. 96, 1982 (1992).
- O. Docenko, M. Tamanis, R. Ferber, E. A. Pazyuk, A. Zaitservkii, A. V. Stolyarov, A. Pashov, H. Knöckel and E. Tiemann, Phys. Rev. A 75, 042503 (2007).
- W. Weickenmeier, U. Diemer, M. Wahl, M. Raab, W. Demtröder, and W. Müller, J. Chem. Phys. 82, 5354 (1985).
- 37. C. Amiot and O. Dulieu, J. Chem. Phys. 117, 5155 (2002).
- B. W. Shore, K. Bergmann, J. Oreg, and S. Rosenwaks, Phys. Rev. A 44, 7442 (1991).
- E. Kuznetsova, P. Pellegrini, R. Côté, M.D. Lukin, and S.F. Yelin, arXiv:0806.0821 (2008).
- M. Greiner, O. Mandel, T. Esslinger, T.W. Hänsch, and I. Bloch, Nature 415, 39 (2002).
- D. Jaksch, V. Venturi, J. I. Cirac, C. J. Williams, and P. Zoller, Phys. Rev. Lett. 89, 040402 (2002).

Precision molecular spectroscopy for ground state transfer of molecular quantum gases

Johann G. Danzl,^a Manfred J. Mark,^a Elmar Haller,^a Mattias Gustavsson,^a

Nadia Bouloufa,^b Olivier Dulieu,^b Helmut Ritsch,^c Russell Hart,^a

and Hanns-Christoph Nägerl^a

^a Institut für Experimentalphysik und Zentrum für Quantenphysik, Universität Innsbruck, Technikerstraße 25, A–6020 Innsbruck, Austria. E-mail: johann.danzl@uibk.ac.at

^b Laboratoire Aimé Cotton, CNRS, Université Paris-Sud, Bât. 505, 91405 Orsay Cedex, France.

^c Institut für Theoretische Physik und Zentrum für Quantenphysik, Universität Innsbruck, Technikerstraße 25, A–6020 Innsbruck, Austria.

December 2, 2008

One possibility for the creation of ultracold, high-phase-space-density quantum gases of molecules in the rovibronic ground state relies on first associating weakly-bound molecules from quantum-degenerate atomic gases on a Feshbach resonance and then transfering the molecules via several steps of coherent two-photon stimulated Raman adiabatic passage (STIRAP) into the rovibronic ground state. Here, in ultracold samples of Cs₂ Feshbach molecules produced out of ultracold samples of Cs atoms, we observe several optical transitions to deeply bound rovibrational levels of the excited 0_u^+ molecular potentials with high resolution. At least one of these transitions, although rather weak, allows efficient STIRAP transfer into the deeply bound vibrational level |v = 73 > of the singlet $X^1 \Sigma_g^+$ ground state potential, as recently demonstrated [1]. From this level, the rovibrational ground state $|v=0, J=0\rangle$ can be reached with one more transfer step. In total, our results show that coherent ground state transfer for Cs₂ is possible using a maximum of two successive two-photon STIRAP processes or one single four-photon STIRAP process.

1 Introduction

Ultracold and dense molecular samples in specific deeply bound rovibrational levels are of high interest for fundamental studies in physics and chemistry. They are expected to find applications in high resolution spectroscopy and fundamental tests [2, 3], few-body collisional physics [4, 5], ultracold chemistry [6], quantum processing [7], and in the field of dipolar quantum gases and dipolar Bose-Einstein condensation [8, 9]. Ideally, full control over the molecular wave function is desired, i.e. full (quantum) control over the internal and external degrees of freedom of the molecules. High phase space densities are needed for molecular quantum gas studies. For many of the envisaged studies and applications, initial preparation of the molecular sample in the rovibronic ground state, i.e. the lowest energy level of the electronic ground state, is desired. Only in this state one can expect sufficient collisional stability.

But how is it possible to produce dense samples of ultracold molecules in the rovibrational ground state? Laser cooling of atoms, which has lead to the production of quantum degenerate atomic Bose and Fermi gases [10], can so far not be adapted to the case of molecular systems as suitable cycling transitions are not available. Versatile non-optical cooling and slowing techniques such as buffer gas cooling and Stark deceleration in combination with molecule trapping [11, 12, 13] have been developed, but high molecular densities and in particular high phase space densities are vet to be reached. An alternative route to producing ultracold molecular samples is given by first producing ultracold atomic samples and then associating molecules out of the atomic sample. While this technique is so far limited to the production of selected species of dimer molecules, it has the advantage that ultra-low temperatures and high particle densities are easily inherited from the atomic precursor sample. There are essentially two association techniques, photoassociation [14] and magnetically induced Feshbach association [15, 16]. In photoassociation experiments [17, 18, 19, 20], ultracold samples of deeply bound molecules have been created. Additional techniques such as vibrational cooling [19] should allow selective pumping into the rovibrational ground state and open up the prospect for high molecular phase space densities. In Feshbach association experiments [21, 22], high-density samples of weakly bound molecules are produced. For dimer molecules composed of Fermions, collisional stability of the highly excited molecules is assured as a result of a Pauli blocking effect, and molecular Bose-Einstein condensation could be achieved in the limit of extremely weak binding [23].

Here, we are interested in combining the techniques of Feshbach association and coherent molecular state transfer to produce quantum gases of molecules in the rovibrational ground state |v=0, J=0> of the lowest electronic state. As usual, v and J are the vibrational and rotational quantum numbers, respectively. The molecules, produced on a Feshbach resonance and hence initially very loosely bound, are to be transferred in a few successive steps of coherent two-photon laser transfer to the rovibrational ground state, acquiring more and more binding energy in each step. The general idea is sketched in Fig.1A for the case of Cs_2 . Each two-photon step involves an excited state level. Population transfer into this level needs to be avoided to prevent loss due to spontaneous emission. One possibility is to use the technique of stimulated Raman adiabatic passage (STIRAP) [24], which is very robust and largely insensitive to laser intensity fluctuations. The scheme has several advantages. First, production of Feshbach molecules out of a quantum degenerate atomic sample can be very efficient [25]. Second, the optical transition rate on the first transition starting from the Feshbach molecules is greatly enhanced in comparison to the free atom case. Further, the scheme is fully coherent, not relying on spontaneous processes, allowing high state selectivity, and involving only a comparatively small number of intermediate levels. A ground state binding energy of typically 0.5 eV for an alkali dimer can be removed essentially without heating the molecular sample, as the differential photon recoil using pairwise co-propagating laser beams driving the two-photon transitions is very small. If losses and off-resonant excitations can be avoided, the scheme essentially preserves phase space density and coherence of the initial particle wave function, allowing the molecular sample to inherit the high initial phase space density from the atomic precursor sample.

Certainly, several challenges have to be met: Going from weakly bound Feshbach to tightly bound ground state molecules corresponds to a large reduction in internuclear distance. Consequently, the radial wave function overlap between successive levels is small, and a compromise has to be found between the number of transitions and the minimum tolerable wave function overlap. To keep the complexity of the scheme low, one or at most two two-photon transitions are desirable. Accordingly, suitable intermediate levels have to be identified that allow a balanced division of wave function overlap, as given by the Franck-Condon factors, between the different transitions. For example, for a four-photon transition scheme with Cs₂ as shown in Fig.1A the Franck-Condon factors are all on the order of 10^{-6} . We emphasize that the identification of the first excited level and hence of the first transition starting from the Feshbach molecules is of crucial importance. Detailed calculations determining the wave function overlap are generally missing, and estimates on the Franck-Condon factors using hypothetical last bound states of either the singlet or triplet potentials of an alkali dimer molecule do not necessarily reflect the transition dipole moments adequately. In addition, for electronic molecular states or energy regions where spectroscopic data is missing, the precise energy of the excited state levels above the atomic threshold is known only with a large uncertainty, which can approach the vibrational spacing of up to a few nanometers. Hence, considerable time has to be spent on searching for weak transitions starting from the initial Feshbach molecules.

In a pioneering experiment, Winkler et al. [26] demonstrated that the STIRAP technique can efficiently be implemented with quantum gases of weakly bound Feshbach molecules. In this work, the transferred molecules, in this case Rb_2 , were still weakly bound with a binding energy of less than 10^{-4} of the binding energy of the rovibronic ground state, and the intermediate excited state level was close to the excited-atom asymptote. Here, we observe several optical transitions starting from a weakly bound Feshbach level to deeply bound rovibrational levels of the mixed excited $(A^1\Sigma_u^+ - b^3\Pi_u) 0_u^+$ molecular potentials of the Cs₂ molecule in a wavelength range from 1118 to 1134 nm, far to the red of the atomic D_1 and D_2 transitions. The Cs_2 molecular potentials are shown in Fig.1A. We observe the levels as loss from an ultracold sample of Cs₂ Feshbach molecules as shown in Fig.1B. We observe two progressions, one that we attribute to the $(A^1\Sigma_u^+ - b^3\Pi_u) 0_u^+$ potentials and one that we associate to the triplet $(1)^{3}\Sigma_{q}^{+}$ potential. From the loss measurements, we determine the transition strengths and find that the stronger transitions should be suitable for STIRAP to an intermediate, deeply bound rovibrational level of the singlet $X^1\Sigma_g^+$ potential with v = 73. Recently, we could implement STIRAP into |v = 73, J = 2 > [1]. For the case of the dimer molecule KRb, Ni et al. [27] could demonstrate quantum gas transfer all the way into the rovibrational ground state |v = 0, J = 0 of the singlet $X^{1}\Sigma^{+}$ molecular potential. Here, the transfer could be achieved in only a single step as a result of the favorable run of the excited state potentials, which is generally the case for heteronuclear molecules composed of alkali atoms [28]. Also recently, transfer to the rovibrational ground state level of the lowest triplet $a^{3}\Sigma_{u}^{+}$ state of Rb₂ could be achieved [29].

2 Preparation of a sample of weakly bound Feshbach molecules

We produce ultracold samples of molecules on two different Feshbach resonances, one near 1.98 mT and one near 4.79 mT [30]. In both cases, essentially following the procedure detailed in Ref.[31], we first produce an ultracold sample of typically 2×10^5 Cs atoms in the lowest hyperfine sublevel F=3, $m_F=3$ in a crossed optical dipole trap. As usual, F is the atomic angular momentum quantum number, and m_F its projection on the magnetic field axis. The trapping light at 1064.5 nm is derived from a singlefrequency, highly-stable Nd:YAG laser. The offset magnetic field value for evaporative cooling is 2.1 mT. We support optical trapping by magnetic levitation with a magnetic field gradient of 3.1 mT/cm. We then produce weakly bound Feshbach molecules out of the atomic sample [22]. We produce a sample every 8 s, i.e. our spectroscopic measurements are performed at a rate of one data point every 8 s. In order to be able to search for optical transitions over large frequency ranges it is advantageous to work with the shortest possible sample preparation times. For this reason we stop evaporative cooling slightly before the onset of Bose-Einstein condensation (BEC), which also makes sample preparation somewhat less critical. The temperature of the initial atomic sample is then typically about 100 nK. At higher temperatures and hence lower phase space densities the molecule production efficiency is reduced, so that there is a trade off between ease of operation and molecule number. We note that for our ground state transfer experiments reported in Ref. [1] we produce a pure atomic BEC at the expense of longer sample preparation times.

The spectrum of weakly-bound Feshbach levels near the two-free-atom asymptote is shown in Fig.2 [30]. For molecule production at the Feshbach resonance at 4.79 mT, we first ramp the magnetic field from the BEC production value to 4.9 mT, about 0.1 mT above the Feshbach resonance. We produce the molecular sample on a downward sweep at a typical sweep rate of 0.025 T/s. The resulting ultracold sample contains up to 11000 molecules, immersed in the bath of the remaining ultracold atoms. The resonance at 4.79 mT is a *d*-wave resonance [30], and hence the molecules are initially of d-wave character, i.e. $\ell = 2$, where ℓ is the quantum number associated with the mechanical rotation of the nuclei. However, there is a weakly bound s-wave Feshbach state ($|s\rangle = |\ell = 0\rangle$) belonging to the open scattering channel right below threshold. This state couples quite strongly to the initial d-wave state, resulting in an avoided state crossing (as shown in the inset to Fig.2), on which the molecules are transferred to the s-wave state $|s\rangle$ upon lowering the magnetic field [30, 1]. Upon further lowering the magnetic field to less than 2.0 mT, the molecules acquire more and more character of a closed channel s-wave state on a second, very broad avoided crossing. Here, we perform spectroscopy in this transition range from open channel to closed channel swave character. At a magnetic field value of 2.0 mT, the binding energy of the molecules is near 5 MHz×h with respect to the $F = 3, m_F = 3$ two-atom asymptote, where h is Planck's constant.

For molecule production at the Feshbach resonance at 1.98 mT, we simply ramp the magnetic field down from the initial BEC production value. Again, we produce an ultracold molecular sample with about 11000 molecules. The molecules in $|g\rangle$ have g-wave character, i.e. $\ell=4$. When we lower the magnetic field to 1.6 mT, the binding energy of the molecules is also near 5 MHz×h with respect to the $F = 3, m_F = 3$ two-atom asymptote. For spectroscopy, we release the molecules from the trap after magnetic field ramping is completed and perform all subsequent experiments in free flight without any other light fields on except for the spectroscopy laser.

For molecule detection in both cases, we reverse the magnetic field ramps [22]. The g-wave molecules are dissociated on the g-wave Feshbach resonance at 1.98 mT, and the s-wave molecules are dissociated on the d-wave Feshbach resonance at 4.79 mT. Prior to the reverse magnetic field ramp, we apply a magnetic field gradient of 3.1 mT/cm for about 5 ms to separate the molecular sample from the atomic sample in a Stern-Gerlach-type experiment. Finally, we detect atoms by standard absorption imaging. The minimum number of molecules that we can detect is on the order of 200 molecules.

3 Spectroscopy

We perform optical spectroscopy on Feshbach molecules in the wavelength region around 1125 nm. Based on selection rules, there are two sets of electronically excited states that we address in the spectroscopic measurements presented here, namely the $(A^1\Sigma_u^+ - b^3\Pi_u) 0_u^+$ coupled state system and the purely triplet $(1)^3\Sigma_g^+$ state. We first discuss transitions to the 0_u^+ coupled state system. Transitions to the latter state are discussed in Sec. 3.2.

3.1 Transitions to the $(\mathbf{A}^1 \Sigma_u^+ - \mathbf{b}^3 \Pi_u) \ 0_u^+$ coupled electronically excited states

For ground state transfer, we are primarily interested in transitions from Feshbach levels to rovibrational levels of the $(A^1\Sigma_u^+ - b^3\Pi_u) 0_u^+$ electronically excited states. In the heavy alkali dimers, most notably in Cs₂, the $A^1\Sigma_u^+$ state and the $b^3\Pi_u$ state are strongly coupled by resonant spin-orbit interaction [32, 33], yielding the 0_u^+ coupled states in Hund's case (c) notation. The singlet component of the 0^+_u states allows us to efficiently couple to deeply bound $X^1\Sigma_q^+$ state levels, specifically to the |v=73, J=2>level of the ground state potential, as has recently been shown in a coherent transfer experiment [1]. We have chosen to do spectroscopy in the wavelength range of 1118 nm to 1134 nm above the $6S_{\frac{1}{2}}+6S_{\frac{1}{2}}$ dissociation threshold of the Cs_2 dimer. This corresponds to a detuning of roughly 2300 $\rm cm^{-1}$ from the cesium D₁ line and to an energy range of approximately 12572 cm^{-1} to 12450 cm^{-1} above the rovibronic ground state $X^1\Sigma_q^+ | v = 0, J = 0 >$. This region was chosen in order to give a balanced distribution of transition strengths in a 4-photon transfer scheme to the rovibronic ground state. In addition, the wavelengths of the four lasers used in the transfer experiments were chosen such that they lie within the wavelength range covered by the infrared fiber-based frequency comb that we use as a frequency reference in the state transfer experiments.

The transitions of interest here lie outside the energy regions for which Fourier transform spectroscopic data was obtained at Laboratoire Aimé Cotton from transitions to the $X^1\Sigma_g^+$ state [34]. The vibrational progression of the 0_u^+ states is highly perturbed by the resonant spin-orbit coupling and exhibits an irregular vibrational spacing. Molecular structure calculations are complicated by the spin-orbit coupling and calculated term values are highly sensitive to the coupling. Prior to the experiments discussed here the absolute energies of the vibrational levels of the $(A^1\Sigma_u^+-b^3\Pi_u) 0_u^+$

excited state levels were poorly known in the region of interest from 1118 nm to 1134 nm. We therefore perform a broad range search by irradiating the weakly-bound Feshbach molecules at a fixed wavelength for a certain irradiation time τ of up to $\tau = 6$ ms and by recording the number of remaining molecules as a function of laser frequency. In one run of the experiment one particular laser frequency is queried. We thus take data points at the repetition rate of our experiment, which is given by the sample preparation time of 8 seconds. Based on the available laser intensity from L_1 and an estimate of the dipole transition moments for the strongest expected lines, we chose a frequency step size of about 100 MHz to 150 MHz for initial line searching. We obtain the laser light at 1118 nm to 1134 nm from a grating-stabilized external cavity diode laser. For coarse frequency scanning, the laser is free running and tuned via a piezoelectric element on the grating of the laser. For more precise measurements, we lock the laser to a narrow-band optical resonator that can be tuned via a piezoelectric element. Fig.3 A shows a typical loss spectrum starting from Feshbach state $|s\rangle$ for excitation near 1126 nm, measured at a magnetic field of 1.98 mT. In this particular case we find three resonances, which we associate with the rotational splitting of the excited state level, J = 5, 3, 1, where J is the rotational quantum number. Based on molecular structure calculations we identify this level as the 225th one of the 0^+_{μ} progression with an uncertainty of about two in the absolute numbering. We zoom in on these three transitions in Fig.3 B. C. and D and record loss resonances at reduced laser intensity in order to avoid saturation of the lines. For these measurements, the laser is locked to the narrow-band optical resonator and the resonator in turn is stabilized to the optical frequency comb to assure reproducibility and long term frequency stability. As one can expect, the loss is strongest on the transition to the $|J=1\rangle$ level, and it is weakest on the transition to |J=5>. The width of all lines gives an excited state spontaneous decay rate of around $2\pi \times 2$ MHz, in agreement with the typical expected lifetimes of excited molecular levels. The transition to |J = 1 > shown in Fig.3 D is of special interest to the current work. It has been used as intermediate excited state level for coherent transfer to $X^1 \Sigma_q^+ | v = 73, J = 2 >$ in our recent experiments [1].

By fitting to a series of such measurements, obtained with different laser intensities, a two level model that takes into account decay from the upper level, we determine the transition strength as given by the normalized Rabi frequency. As a result of optical excitation, for small saturation the number N of Feshbach molecules decays as a function of laser detuning Δ_1 according to $N(\Delta_1) = N_0 \exp(-\tau \Omega_1^2/(\Gamma(1 + 4\pi^2 \Delta_1^2/\Gamma^2))))$, where N_0 is the molecule number without laser irradiation and τ is the irradiation time. From the fit we obtain the Rabi frequency on resonance Ω_1 and the excited state spontaneous decay rate Γ . We determine the normalized Rabi frequency to $\Omega_1 = 2\pi \times 2 \text{ kHz} \sqrt{I/(\text{mW/cm}^2)}$ for |J=1>, where I is the laser intensity. This value is sufficient to perform STIRAP given the available laser power [1]. The corresponding transition strengths for |J=3> and |J=5> are $\Omega_1 = 2\pi \times 0.3 \text{ kHz} \sqrt{I/(\text{mW/cm}^2)}$ and $\Omega_1 = 2\pi \times 0.1 \text{ kHz} \sqrt{I/(\text{mW/cm}^2)}$, respectively. The absolute values of these transition strengths bear an estimated uncertainty of 20 % because the laser beam parameters for the spectroscopy laser are not well determined.

We also record the time dependence of the molecular loss on some of the stronger lines. For this, we step the laser irradiation time τ from 0 to 150 μ s, while laser L_1 is kept on resonance. The result is shown in Fig.4 A for the transition at 1126.173 nm for two different values of the excitation laser intensity.

We note that the transition strength for a particular line starting from Feshbach level $|s\rangle$ strongly depends on the value of the magnetic field, as evidenced in Fig.4 B. Loss resonances for the transition at 1126.173 nm at 1.9 mT and 2.2 mT are shown. For ground state transfer [1], we choose a magnetic field of around 1.9 mT, which is somewhat below the magnetic field region where state $|s\rangle$ is strongly curved, but above the avoided state crossing with state $|g_2\rangle$, as seen in Fig.2. The pronounced bending of $|s\rangle$ is the result of a strong avoided crossing between two s-wave Feshbach levels [30]. For magnetic field values beyond 3.0 mT the level $|s\rangle$ can be associated to the $F_1=3, F_2=3$ asymptote, where $F_i, i=1, 2$, is the atomic angular momentum quantum number of the *i*-th atom, respectively. Below 2.0 mT the level $|s\rangle$ can be associated to the $F_1 = 4, F_2 = 4$ asymptote. It is hence of closed channel character and much more deeply bound with respect to its potential asymptote, effectively by twice the atomic hyperfine splitting, improving the radial wave function overlap with the excited state levels. This increases the transition strength. Trivially, the resonance frequency is shifted as the binding energy is reduced for larger magnetic field values. Coupling to the excited state level is reduced from $\Omega_1 = 2\pi \times 2 \text{ kHz} \sqrt{I/(\text{mW/cm}^2)}$ to $\Omega_1 = 2\pi \times 1$

kHz $\sqrt{I/(\text{mW/cm}^2)}$ when the magnetic field is changed from 1.9 mT to 2.2 mT.

As will be discussed in Sec.4 it is advantageous to be able to choose different Feshbach states as a starting state for ground state transfer experiments. Therefore, we probe transitions from Feshbach level $|g\rangle$ to $(A^1\Sigma_u^+ - b^3\Pi_u) 0_u^+$ levels. Fig.5 shows loss resonances to the same excited state levels as shown in Fig.3, only that now the initial Feshbach level is $|g\rangle$ instead of $|s\rangle$. In this case, the transition to $|J = 3\rangle$ is the strongest, while the transition to $|J = 1\rangle$ is very weak, but can be detected. A comparison of the transition strengths from $|g\rangle$ to the excited state level $|J = 3\rangle$, giving $\Omega_1 = 2\pi \times 1$ kHz $\sqrt{I/(\text{mW/cm}^2)}$ versus $|s\rangle$ to $|J = 1\rangle$ giving $\Omega_1 = 2\pi \times 2$ kHz $\sqrt{I/(\text{mW/cm}^2)}$ shows that level $|g\rangle$ could also be potentially used as a starting level for coherent population transfer to deeply bound levels of the ground state but requires

for coherent population transfer to deeply bound levels of the ground state but requires longer STIRAP times in order to assure sufficient adiabaticity [24]. The $|J=3\rangle$ excited state level in turn couples to $|J=2\rangle$ in the ground state, the level used in our previous work [1].

In addition to the transition near 1126 nm we find a series of other excited state levels that we assign to the $(A^1\Sigma_u^+ - b^3\Pi_u) 0_u^+$ coupled state system. These are listed in Table 1. The assignment to either the $(A^1\Sigma_u^+ - b^3\Pi_u) 0^+_u$ system or to the $(1)^3\Sigma_q^+$ electronically excited state discussed below is primarily based on the spacing between neighboring vibrational levels and in addition on the pattern of loss resonances associated with each particular vibrational level. Resonant spin-orbit coupling in the case of the 0_u^+ states leads to an irregular vibrational spacing. In contrast, the $(1)^3 \Sigma_g^+$ state is not perturbed by spin-orbit interaction and therefore has a regular vibrational progression. The levels near 1126 nm and near 1123 nm have been used to detect dark resonances with deeply bound levels of the $X^1\Sigma_q^+$ state [1]. The ability to couple to these essentially purely singlet ground state levels unambiguously assigns the corresponding excited state levels to the 0^+_u system. The data given in Table 1 does not represent a fully exhaustive study of the $(A^1\Sigma^+_u - b^3\Pi_u) 0^+_u$ coupled states in the wavelength range of interest. In fact, for the most part we observe those levels of the 0^+_{μ} system that have a dominant $A^1\Sigma^+_u$ state contribution, as determined from molecular structure calculations.

3.2 Transitions to the $(1)^{3}\Sigma_{q}^{+}$ electronically excited state

The Feshbach levels that serve as starting levels for the spectroscopy are of mixed $X^1\Sigma_g^+$ and $a^3\Sigma_u^+$ character. In the wavelength range explored here, excitation to the $(1)^3\Sigma_g^+$ electronically excited triplet state is possible from the $a^{3}\Sigma_{u}^{+}$ component of the Feshbach molecules. In fact, for a heavy molecule as Cs₂, the $(1)^{3}\Sigma_{g}^{+}$ state is better described by the two separate electronic states 0_g^- and 1_g , denoted by Hund's case (c) notation. The $(1)^{3}\Sigma_{q}^{+}$ has been previously studied by Fourier transform spectroscopy [35]. This state is not of prime interest for the present work as transitions from this state down to the $X^1\Sigma_q^+$ ground state are expected to be strongly suppressed, but would be important for STIRAP transfer into the rovibrational ground state level of the shallow triplet $a^{3}\Sigma_{n}^{+}$ potential [29]. Certainly, it is important to be able to distinguish rovibrational levels belonging to the $(1)^{3}\Sigma_{q}^{+}$ state from the ones belonging to the 0_{u}^{+} system, because otherwise time would be wasted in searching for ground state dark resonances that are very weak or even do not exist. Fig.6 A shows a typical loss spectrum for one of the lines that we detected near 1127.37 nm. Due to hyperfine splitting, levels of triplet character exhibit a much richer substructure than the 0^+_{μ} levels used for ground state transfer. Several components can be identified as a result of rotational and excited state hyperfine splitting. Zoomed-in regions are shown in Fig.6 B, C, D, and E. We have observed a regularly spaced series of optical transitions which we attribute to the $(1)^{3}\Sigma_{q}^{+}$ excited state as listed in Table 1. The level energies are well reproduced by the Dunham coefficients determined in Ref.[35]. The vibrational numbering used here is the same as in that work. However, it relies on the absolute energy position of the potential, T_e , which was not determined precisely in Ref. [35]. By fixing T_e to the value given in Ref. [35] we get good agreement with our data.

4 Conclusion

We have performed optical spectroscopy starting from weakly bound Cs₂ Feshbach molecules into deeply bound rovibrational levels of the mixed excited state 0^+_u system and the excited triplet $(1)^{3}\Sigma_{g}^{+}$ state. At least one of the observed transitions, namely the one at 1126.173 nm starting from the Feshbach level $|s\rangle$ to the excited level |v'=225, J=1> of the 0_u^+ system, at an offset magnetic field value of 1.9 mT, is strong enough to allow efficient STIRAP transfer into deeply bound rovibrational levels of the singlet $X^1\Sigma_q^+$ ground state potential. The use of this transition for STIRAP has recently been demonstrated in Ref.[1]. In that work, the deeply bound rovibrational level |v=73, J=2> of the $X^1\Sigma_g^+$ ground state potential was populated in the molecular quantum gas regime with 80% efficiency. The rovibrational ground state |v=0, J=0>of the $X^1\Sigma_a^+$ ground state potential can thus be reached from the atomic threshold with a maximum of two two-photon STIRAP transfers. Dark resonances connecting |v=73, J=2> to |v=0, J=0> have recently been observed [36], and two-step STIRAP into |v=0, J=0> has recently been implemented [37]. For future experiments, the use of Feshbach level $|g\rangle$ as the initial state might be advantageous. Level $|g\rangle$ can be more easily populated, as the Feshbach resonance connected to this level is at a low magnetic field value of 1.98 mT [30], where the atomic background scattering length has a moderate value of 155 a_0 , where a_0 is Bohr's radius. The use of this resonance avoids excitation of collective motion of the atomic BEC as a result of a large mean field interaction near the Feshbach resonance at 4.79 mT [1], where the atomic background scattering length is about 935 a₀. The transition starting from level $|g\rangle$ appears to be strong enough to allow STIRAP, this time via the excited state level $|v'=225, J=3\rangle$ of the 0_u^+ system. An attractive strategy for the production of a BEC of ground state molecules relies on the addition of a three-dimensional optical lattice. Starting from the atomic BEC, pairs of atoms at individual lattice sites can be produced in a superfluidto-Mott-insulator transition [38] with high efficiencies of almost 50% [39]. These pairs can then be very efficiently associated on a Feshbach resonance [40] and subsequently transferred to the rovibronic ground state with STIRAP. The lattice has the advantage of shielding the molecules against inelastic collisions during the association process and subsequent state transfer. In particular, it should allow long STIRAP pulse durations, allowing us to resolve the weak hyperfine structure of ground state molecules [41]. As proposed by Jaksch *et al.* [42], dynamical melting of the Mott-insulator state should ideally result in the formation of a BEC of molecules in the rovibronic ground state in a Mott-insulator-to-superfluid-type transition.

5 Acknowledgements

We are indebted to R. Grimm for generous support and we thank T. Bergeman, H. Salami, J. Hutson, J. Aldegunde, and E. Tiemann for valuable discussions. We gratefully acknowledge funding by the Austrian Ministry of Science and Research (BMWF) and the Austrian Science Fund (FWF) in form of a START prize grant and by the European Science Foundation (ESF) in the framework of the EuroQUAM collective research project QuDipMol. R.H. acknowledges support by the European Union in form of a Marie-Curie International Incoming Fellowship (IIF).

References

- J. G. Danzl, E. Haller, M. Gustavsson, M. J. Mark, R. Hart, N. Bouloufa, O. Dulieu, H. Ritsch, and H.-C. Nägerl, *Science*, 2008, **321**, 1062, published online 10 July 2008; 10.1126/science.1159909.
- [2] T. Zelevinsky, S. Kotochigova, and J. Ye, Phys. Rev. Lett., 2008, 100, 043201.
- [3] D. DeMille, S. Sainis, J. Sage, T. Bergeman, S. Kotochigova, and E. Tiesinga, *Phys. Rev. Lett.*, 2008, **100**, 043202.
- [4] P. Staanum, S. D. Kraft, J. Lange, R. Wester, and M. Weidemüller, *Phys. Rev. Lett.*, 2006, **96**, 023201.
- [5] N. Zahzam, T. Vogt, M. Mudrich, D. Comparat, and P. Pillet, *Phys. Rev. Lett.*, 2006, **96**, 023202.
- [6] R. V. Krems, Int. Rev. Phys. Chem., 2005, 24, 99.
- [7] D. DeMille, Phys. Rev. Lett., 2002, 88, 067901.
- [8] K. Góral, L. Santos, and M. Lewenstein, Phys. Rev. Lett., 2002, 88, 170406.
- [9] M. Baranov, L. Dobrek, K. Góral, L. Santos, and M. Lewenstein, *Phys. Scr.*, 2002, T102, 74.
- [10] K. Southwell (ed.), Ultracold matter, Nature (Insight), 2002, 416, 205.
- [11] J. Doyle, B. Friedrich, R.V. Krems, and F. Masnou-Seeuws, *Eur. Phys. J. D*, 2004, 31, 149.
- [12] R. V. Krems, Phys. Chem. Chem. Phys., 2008, 10, 4079.
- [13] S. Y. T. van de Meerakker, H. L. Bethlem, and G. Meijer, Nature Phys., 2008, 4, 595.
- [14] For a review on photoassociation, see: K. M. Jones, E. Tiesinga, P. D. Lett, and P. S. Julienne, *Rev. Mod. Phys.*, 2006 78, 483.
- [15] T. Köhler, K. Góral, and P. S. Julienne, Rev. Mod. Phys., 2006, 78, 1311.
- [16] F. Ferlaino, S. Knoop, and R. Grimm, in: Cold Molecules: Theory, Experiment, Applications, edt. by R. V. Krems, B. Friedrich, and W. C. Stwalley (publication expected in March 2009), preprint at arXiv:0809.3920.
- [17] A. N. Nikolov, J. R. Ensher, E. E. Eyler, H. Wang, W. C. Stwalley, and P. L. Gould, *Phys. Rev. Lett.*, 2000, 84, 246.
- [18] J. M. Sage, S. Sainis, T. Bergeman, and D. DeMille, Phys. Rev. Lett., 2005, 94, 203001.
- [19] M. Viteau, A. Chotia, M. Allegrini, N. Bouloufa, O. Dulieu, D. Comparat, and P. Pillet, *Science*, 2008, **321**, 232.
- [20] J. Deiglmayr, A. Grochola, M. Repp, K. Mörtlbauer, C. Glück, J. Lange, O. Dulieu, R. Wester, and M. Weidemüller, *Phys. Rev. Lett.*, 2008, **101**, 133004.

- [21] C. A. Regal, C. Ticknor, J. L. Bohn, and D. S. Jin, Nature, 2003, 424, 47.
- [22] J. Herbig, T. Kraemer, M. Mark, T. Weber, C. Chin, H.-C. Nägerl, and R. Grimm, *Science*, 2003, **301**, 1510, published online 21 August 2003; 10.1126/science.1088876.
- [23] For an overview, see: Ultracold Fermi Gases, Proceedings of the International School of Physics Enrico Fermi, Course CLXIV, edited by M. Inguscio, W. Ketterle, and C. Salomon (IOS Press, Amsterdam, 2008).
- [24] K. Bergmann, H. Theuer and B. W. Shore, Rev. Mod. Phys., 1998, 70, 1003.
- [25] M. Mark, T. Kraemer, J. Herbig, C. Chin, H.-C. Nägerl, and R. Grimm, *Europhys. Lett.*, 2005, **69**, 706.
- [26] K. Winkler, F. Lang, G. Thalhammer, P. v. d. Straten, R. Grimm, and J. Hecker Denschlag, *Phys. Rev. Lett.*, 2007, 98, 043201.
- [27] K.-K. Ni, S. Ospelkaus, M. H. G. de Miranda, A. Pe'er, B. Neyenhuis, J. J. Zirbel, S. Kotochigova, P. S. Julienne, D. S. Jin, and J. Ye, published online September 18, 2008, *Science* DOI: 10.1126/science.1163861.
- [28] W. C. Stwalley, Eur. Phys. J. D, 2004, 31, 221.
- [29] F. Lang, K. Winkler, C. Strauss, R. Grimm, and J. Hecker Denschlag, Phys. Rev. Lett., 2008, 101, 133005.
- [30] M. Mark, F. Ferlaino, S. Knoop, J. G. Danzl, T. Kraemer, C. Chin, H.-C. Nägerl, and R. Grimm, *Phys. Rev. A*, 2007, **76**, 042514.
- [31] T. Weber, J. Herbig, M. Mark, H.-C. Nägerl, and R. Grimm, *Science*, 2003, 299, 232.
- [32] O. Dulieu and P. Julienne, J. Chem. Phys., 1995, 103, 60.
- [33] C. Amiot, O. Dulieu, and J. Vergès, Phys. Rev. Lett., 1999, 83, 2316.
- [34] H. Salami, T. Bergeman, O. Dulieu, D. Li, F. Xie, and L. Li, manuscript in preparation (2008).
- [35] C. Amiot and J. Vergès, Chem. Phys. Lett., 1985, 116, 273.
- [36] M. J. Mark, J. G. Danzl, E. Haller, M. Gustavsson, N. Bouloufa, O. Dulieu, H. Salami, T. Bergeman, H. Ritsch, R. Hart, and H.-C. Nägerl, manuscript submitted for publication, 2008, preprint at arXiv:0811.0695.
- [37] J. G. Danzl, M. J. Mark, E. Haller, M. Gustavsson, N. Bouloufa, O. Dulieu, H. Salami, T. Bergeman, H. Ritsch, R. Hart, H.-C. Nägerl, manuscript in preparation.
- [38] M. Greiner, O. Mandel, T. Esslinger, T. W. Hänsch, and I. Bloch, Nature, 2002, 415, 39.
- [39] S. Dürr, 2008, private communication.
- [40] G. Thalhammer, K. Winkler, F. Lang, S. Schmid, R. Grimm, and J. Hecker Denschlag, *Phys. Rev. Lett.*, 2006, 96, 050402.

- [41] J. Aldegunde and J. M. Hutson, 2008, preprint at arXiv:0810.4709.
- [42] D. Jaksch, V. Venturi, J. I. Cirac, C. J. Williams, and P. Zoller, *Phys. Rev. Lett.*, 2002, 89, 040402.

Table 1: Observed excited state levels in the wavelength range from 1118 nm to 1134 nm. Transitions were measured from Feshbach state $|s\rangle$ to the first electronically excited state, addressing both $(A^1 \Sigma_u^+ - b^3 \Pi_u) 0_u^+$ levels and $(1)^3 \Sigma_q^+$ levels. Levels are given according to the excitation wavelength (WL) from $|s\rangle$, which essentially corresponds to the $F=3, m_F=3$ two-atom asymptote. The data is taken at a magnetic field of 1.98 mT. Wavemeter accuracy is about 0.001 nm. The energy of these levels above the rovibronic ground state $X^1 \Sigma_g^+ | v = 0, J = 0 >$ is given in the second column, where the binding energy of the rovibronic ground state is taken from Ref.[1]. The assignment to either the coupled $(A^1\Sigma_u^+ - b^3\Pi_u)0_u^+$ system or to the $(1)^3\Sigma_q^+$ is based on the vibrational spacing and similarities in the substructure of the levels. The levels marked with * have been used for dark resonance spectroscopy coupling to deeply bound levels of the $X^1\Sigma_q^+$ state [1]. The ability to couple to such levels unambiguously reflects an important singlet component stemming from the $A^1\Sigma^+_u$ state and therefore clearly assigns these levels to the 0^+_{μ} system. The quantum numbers given for the 0^+_{μ} levels are coupled channels quantum numbers derived from molecular structure calculations and bear an uncertainty of two in the absolute numbering. The calculations show that these levels have about 70% $A^1 \Sigma_u^+$ state contribution. Two further levels observed near 1120.17 nm and 1117.16 nm that belong to the 0_u^+ progression are not given in the table since no further measurements have been done on these levels. The level near 1129.5 nm exhibits a somewhat richer structure than the other levels assigned to 0^+_{μ} and than exemplified in Fig. 3. Levels assigned to the $(1)^{3}\Sigma_{g}^{+}$ state form a regular vibrational progression and show a more complex substructure than the levels attributed to the 0_{μ}^{+} system, as exemplified in Fig. 6. For these levels, the transition wavelength to one of the most prominent features is given, since an in-depth analysis of the rotational and hyperfine structure remains to be done. The vibrational numbering for the $(1)^3 \Sigma_a^+$ levels is the same as in Ref [35].

WL [nm]	Energy above	Assignment
	$X^1\Sigma_a^+ v=0>$	
	$[\mathrm{cm}^{-1}]$	
1132.481	12458.875	$0_u^+ v' = 221, J = 1 >$
1129.492	12482.245	0^{+}_{u}
1126.173*	12508.332	$0_u^+ v' = 225, J = 1 >$
1123.104*	12532.598	$0_u^+ v' = 226, J = 1 >$
1133.680	12449.536	$(1)^{3}\Sigma_{a}^{+} v'=32>$
1130.510	12474.274	$(1)^{3}\Sigma_{a}^{+} v' = 33 >$
1127.379	12498.838	$(1)^{3}\Sigma_{a}^{+} v'=34>$
1124.274	12523.334	$(1)^{3}\Sigma_{q}^{+} v' = 35 >$
1121.196	12547.756	$(1)^{3}\Sigma_{q}^{+} v' = 36 >$
1118.155	12572.013	$(1)^{3}\Sigma_{g}^{+} v'=37>$

- Fig. 1 (\mathbf{A}) Simplified molecular level scheme for Cs_2 showing the relevant ground state and excited state potentials involved in rovibrational ground state transfer. Molecules in a weakly bound Feshbach level $|1\rangle = |v \approx 155\rangle$ (not resolved near the $6S_{\frac{1}{2}} + 6S_{\frac{1}{2}}$ two-atom asymptote, but shown in Fig.2) are to be transferred to the rovibrational ground state level |5>=|v=0, J=0>of the singlet $X^1\Sigma_q^+$ potential with a binding energy of 3629 cm⁻¹ by two sequential two-photon STIRAP processes involving lasers L_1 and L_2 near 1126 nm and 1006 nm and lasers L_3 and L_4 near 1351 nm and 1003 nm. The intermediate ground state level $|3\rangle = |v = 73, J = 2\rangle$ has a binding energy of 1061 cm⁻¹. (**B**) Probing candidate levels for $|2\rangle$ belonging to the electronically excited coupled $(A^1 \Sigma_u^+ - b^3 \Pi_u) 0_u^+$ potentials. Here, we search for $|2\rangle$ in loss spectroscopy with laser L_1 in a region near 8890 $\rm cm^{-1}$ above the 6S $_{\frac{1}{2}}$ + 6S $_{\frac{1}{2}}$ asymptote, corresponding to an excitation wavelength range of 1118 to 1134 nm. The wiggly arrow indicates loss from the excited levels due to spontaneous emission. Also shown is the excited $(1)^{3}\Sigma_{g}^{+}$ potential, for which we find several levels.
- Fig. 2 Initial Feshbach molecule production: Zeeman diagram showing the energy of weakly bound Feshbach levels [30] and the Feshbach resonances (FR) used in the present work. The binding energy is given with respect to the $F = 3, m_F = 3$ two-atom asymptote. The molecules are produced either on a d-wave Feshbach resonance at 4.79 mT (see inset) and then transferred to the weakly bound s-wave state $|s\rangle$ on an avoided state crossing, or on a g-wave Feshbach resonance at 1.98 mT, resulting in molecules in level |q>. In the first case, further lowering of the magnetic offset field to below 2.0 mT changes the character of the $|s\rangle$ level from open-channel to closed-channel dominated [30]. The levels $|s\rangle$ and $|g\rangle$ are both candidate levels for the initial level |1> shown in Fig.1. For completeness, further g-wave Feshbach levels, $|g_1\rangle$, $|g_2\rangle$, and $|g_3\rangle$ are shown. Level $|g_2\rangle$ connects $|g\rangle$ to $|s\rangle$ and can be used for Feshbach state transfer [30]. Level $|g_3\rangle$ is a further interesting candidate level for $|1\rangle$ with low nuclear spin contribution [30].
- Fig. 3 Loss resonances for excitation from the initial Feshbach level $|s\rangle$ to the 0_u^+ system. (A) Typical scan showing the relative number of molecules in $|s\rangle$ as a function of laser wavelength λ_1 near 1126 nm. Three resonances can be identified, corresponding to |J=5>, |J=3>, and |J=1>, from left to right. The sample is irradiated with laser light at an intensity of 1×10^6 mW/cm² for $\tau = 200 \ \mu s$. The laser is locked to a narrow band optical resonator that is tuned via a piezoelectric element with a step size of approximately 40 MHz. Wavelength is measured on a homebuilt wavemeter. The molecule number is normalized to the atom number measured in the same individual realization of the experiment to cancel out fluctuations that stem from shot-to-shot atom number fluctuations and then normalized to unity. (\mathbf{B}) , (\mathbf{C}) , and (\mathbf{D}) show measurements of the three individual lines with |J=5>, |J=3>, and |J=1> at reduced intensity in order to avoid saturation. The solid lines represent fits as described in the text. The spectroscopy laser is stabilized to an optical resonator and the resonator is in turn referenced to an optical frequency
comb, which allows precise and reproducible tuning of the frequency. The transition to $|J=1\rangle$ in panel (**D**) is recorded at an intensity of 1.5×10^4 mW/cm² (circles) and 6×10^3 mW/cm² (triangles), (**B**) and (**C**) are recorded at 1×10^6 mW/cm² and 2×10^5 mW/cm², respectively. The pulse duration is $\tau = 10 \ \mu$ s.

- Fig. 4 Loss of molecules for excitation near 1126.173 nm from Feshbach level |s>. (A) Time dependence of molecular loss on resonance at 1126.173 nm for two different laser intensities, $5.7 \times 10^5 \text{ mW/cm}^2$ (circles) and $2.1 \times 10^5 \text{ mW/cm}^2$ (triangles). The magnetic offset field is 1.9 mT. The fitted exponential decay gives the decay constants $\tau_d = 9.7 \pm 0.6 \ \mu$ s (circles) and $\tau_d = 25.5 \pm 1 \ \mu$ s (triangles). (B) Loss of molecules in $|s\rangle$ as a function of laser detuning Δ_1 near 1126 nm with an irradiation time of $\tau = 10 \ \mu$ s for two values of the magnetic field, 1.9 mT (circles) and 2.2 mT (triangles). In both cases, the excited state spontaneous decay rate was determined to $\approx 2\pi \times 2 \ \text{MHz}$. At higher magnetic fields, Feshbach level $|s\rangle$ acquires more open-channel character, reducing radial wave function overlap with the excited rovibrational levels. The shift in transition frequency is essentially the result of the change in binding energy as seen in Fig. 2.
- Fig. 5 Loss resonances for excitation from the initial Feshbach level |g>. (A), (B), and (C) show the loss for excitation to |J = 5>, |J = 3>, and |J = 1>, corresponding to the resonances shown in Fig.3. The laser intensities are 1.5×10^4 mW/cm² for panel (A) and for the circles in panel (B). The second resonance in (B) (triangles) is measured with 5.6×10^3 mW/cm². (C) The line at 1126.173 nm is measured at 1×10^6 mW/cm². All measurements are done with an irradiation time of $\tau = 10 \ \mu$ s. From a series of such measurements at different intensities we determine the line strengths for |J = 5>, |J = 3>, and |J = 1> to $\Omega_1 = 2\pi \times 1$ kHz $\sqrt{I/(mW/cm^2)}$, $\Omega_1 = 2\pi \times 1$ kHz $\sqrt{I/(mW/cm^2)}$, and $\Omega_1 = 2\pi \times 0.1$ kHz $\sqrt{I/(mW/cm^2)}$, respectively.
- Fig. 6 Loss of molecules for excitation near 1127.17 nm from Feshbach level $|s\rangle$ to the triplet $(1)^{3}\Sigma_{q}^{+}$ state. (A) represents a broad scan with laser irradiation at an intensity of 5×10^5 mW/cm² for $\tau = 100 \ \mu s$ at a step size of 20 MHz. A rich structure due to rotation and excited state hyperfine splitting can be seen which is qualitatively different from the spectrum shown in Fig.3. The lines are greatly broadened by the high intensity and long irradiation time. The spectroscopy laser is locked to a narrow band optical resonator that is stepped via a piezoelectric element. Scans of about 750 MHz were recorded as a function of piezo voltage on the resonator. Voltage was converted to wavelength for each scan by a linear interpolation. (\mathbf{B}) - (\mathbf{E}) represent scans over some of the observed features at a reduced intensity of $8 \times 10^4 \text{ mW/cm}^2$ and an irradiation time of $\tau = 10 \ \mu s$ in order to reduce broadening of the lines. The step size is about 7 MHz. Resonator piezo voltage is converted to frequency with an estimated error of 10 %. The absolute wavelength accuracy is limited by wavemeter calibration to 0.001 nm, the relative accuracy is about a factor of 10 better. The vertical arrows indicate weak lines that have

been verified in additional scans with higher power. In panel (\mathbf{E}) the power was somewhat increased for an additional measurement (triangles) that emphasizes such a weak line.



Figure 1: (A) Simplified molecular level scheme for Cs₂ showing the relevant ground state and excited state potentials involved in rovibrational ground state transfer. Molecules in a weakly bound Feshbach level $|1\rangle = |v \approx 155\rangle$ (not resolved near the $6S_{\frac{1}{2}} + 6S_{\frac{1}{2}}$ two-atom asymptote, but shown in Fig.2) are to be transferred to the rovibrational ground state level $|5\rangle = |v = 0, J = 0\rangle$ of the singlet $X^{1}\Sigma_{g}^{+}$ potential with a binding energy of 3629 cm⁻¹ by two sequential two-photon STIRAP processes involving lasers L_{1} and L_{2} near 1126 nm and 1006 nm and lasers L_{3} and L_{4} near 1351 nm and 1003 nm. The intermediate ground state level $|3\rangle = |v = 73, J = 2\rangle$ has a binding energy of 1061 cm⁻¹. (B) Probing candidate levels for $|2\rangle$ belonging to the electronically excited coupled $(A^{1}\Sigma_{u}^{+}-b^{3}\Pi_{u}) 0_{u}^{+}$ potentials. Here, we search for $|2\rangle$ in loss spectroscopy with laser L_{1} in a region near 8890 cm⁻¹ above the $6S_{\frac{1}{2}} + 6S_{\frac{1}{2}}$ asymptote, corresponding to an excitation wavelength range of 1118 to 1134 nm. The wiggly arrow indicates loss from the excited levels due to spontaneous emission. Also shown is the excited $(1)^{3}\Sigma_{g}^{+}$ potential, for which we find several levels.



Figure 2: Initial Feshbach molecule production: Zeeman diagram showing the energy of weakly bound Feshbach levels [30] and the Feshbach resonances (FR) used in the present work. The binding energy is given with respect to the F=3, $m_F=3$ two-atom asymptote. The molecules are produced either on a *d*-wave Feshbach resonance at 4.79 mT (see inset) and then transferred to the weakly bound *s*-wave state $|s\rangle$ on an avoided state crossing, or on a *g*-wave Feshbach resonance at 1.98 mT, resulting in molecules in level $|g\rangle$. In the first case, further lowering of the magnetic offset field to below 2.0 mT changes the character of the $|s\rangle$ level from open-channel to closed-channel dominated [30]. The levels $|s\rangle$ and $|g\rangle$ are both candidate levels for the initial level $|1\rangle$ shown in Fig.1. For completeness, further *g*-wave Feshbach levels, $|g_1\rangle$, $|g_2\rangle$, and $|g_3\rangle$ are shown. Level $|g_2\rangle$ connects $|g\rangle$ to $|s\rangle$ and can be used for Feshbach state transfer [30]. Level $|g_3\rangle$ is a further interesting candidate level for $|1\rangle$ with low nuclear spin contribution [30].



Figure 3: Loss resonances for excitation from the initial Feshbach level $|s\rangle$ to the 0_u^+ system. (A) Typical scan showing the relative number of molecules in $|s\rangle$ as a function of laser wavelength λ_1 near 1126 nm. Three resonances can be identified, corresponding to $|J = 5\rangle$, $|J = 3\rangle$, and $|J = 1\rangle$, from left to right. The sample is irradiated with laser light at an intensity of $1 \times 10^6 \text{ mW/cm}^2$ for $\tau = 200 \ \mu\text{s}$. The laser is locked to a narrow band optical resonator that is tuned via a piezoelectric element with a step size of approximately 40 MHz. Wavelength is measured on a home-built wavemeter. The molecule number is normalized to the atom number measured in the same individual realization of the experiment to cancel out fluctuations that stem from shot-to-shot atom number fluctuations and then normalized to unity. (\mathbf{B}) , (\mathbf{C}) , and (D) show measurements of the three individual lines with $|J = 5\rangle$, $|J = 3\rangle$, and |J=1> at reduced intensity in order to avoid saturation. The solid lines represent fits as described in the text. The spectroscopy laser is stabilized to an optical resonator and the resonator is in turn referenced to an optical frequency comb, which allows precise and reproducible tuning of the frequency. The transition to $|J=1\rangle$ in panel (**D**) is recorded at an intensity of 1.5×10^4 mW/cm² (circles) and 6×10^3 mW/cm² (triangles), (B) and (C) are recorded at 1×10^6 mW/cm² and 2×10^5 mW/cm², respectively. The pulse duration is $\tau = 10 \ \mu s$.



Figure 4: Loss of molecules for excitation near 1126.173 nm from Feshbach level |s>. (A) Time dependence of molecular loss on resonance at 1126.173 nm for two different laser intensities, $5.7 \times 10^5 \text{ mW/cm}^2$ (circles) and $2.1 \times 10^5 \text{ mW/cm}^2$ (triangles). The magnetic offset field is 1.9 mT. The fitted exponential decay gives the decay constants $\tau_d = 9.7 \pm 0.6 \ \mu\text{s}$ (circles) and $\tau_d = 25.5 \pm 1 \ \mu\text{s}$ (triangles). (B) Loss of molecules in |s> as a function of laser detuning Δ_1 near 1126 nm with an irradiation time of $\tau = 10 \ \mu\text{s}$ for two values of the magnetic field, 1.9 mT (circles) and 2.2 mT (triangles). In both cases, the excited state spontaneous decay rate was determined to $\approx 2\pi \times 2 \text{ MHz}$. At higher magnetic fields, Feshbach level |s> acquires more open-channel character, reducing radial wave function overlap with the excited rovibrational levels. The shift in transition frequency is essentially the result of the change in binding energy as seen in Fig. 2.



Figure 5: Loss resonances for excitation from the initial Feshbach level $|g\rangle$. (A), (B), and (C) show the loss for excitation to $|J=5\rangle$, $|J=3\rangle$, and $|J=1\rangle$, corresponding to the resonances shown in Fig.3. The laser intensities are 1.5×10^4 mW/cm² for panel (A) and for the circles in panel (B). The second resonance in (B) (triangles) is measured with 5.6×10^3 mW/cm². (C) The line at 1126.173 nm is measured at 1×10^6 mW/cm². All measurements are done with an irradiation time of $\tau = 10 \ \mu$ s. From a series of such measurements at different intensities we determine the line strengths for $|J=5\rangle$, $|J=3\rangle$, and $|J=1\rangle$ to $\Omega_1 = 2\pi \times 1$ kHz $\sqrt{I/(\text{mW/cm}^2)}$, $\Omega_1 = 2\pi \times 1$ kHz $\sqrt{I/(\text{mW/cm}^2)}$, and $\Omega_1 = 2\pi \times 0.1$ kHz $\sqrt{I/(\text{mW/cm}^2)}$, respectively.



Figure 6: Loss of molecules for excitation near 1127.17 nm from Feshbach level $|s\rangle$ to the triplet $(1)^3 \Sigma_g^+$ state. (A) represents a broad scan with laser irradiation at an intensity of 5×10^5 mW/cm² for $\tau = 100 \ \mu s$ at a step size of 20 MHz. A rich structure due to rotation and excited state hyperfine splitting can be seen which is qualitatively different from the spectrum shown in Fig.3. The lines are greatly broadened by the high intensity and long irradiation time. The spectroscopy laser is locked to a narrow band optical resonator that is stepped via a piezoelectric element. Scans of about 750 MHz were recorded as a function of piezo voltage on the resonator. Voltage was converted to wavelength for each scan by a linear interpolation. (\mathbf{B}) - (\mathbf{E}) represent scans over some of the observed features at a reduced intensity of $8 \times 10^4 \text{ mW/cm}^2$ and an irradiation time of $\tau = 10 \ \mu s$ in order to reduce broadening of the lines. The step size is about 7 MHz. Resonator piezo voltage is converted to frequency with an estimated error of 10 %. The absolute wavelength accuracy is limited by wavemeter calibration to 0.001nm, the relative accuracy is about a factor of 10 better. The vertical arrows indicate weak lines that have been verified in additional scans with higher power. In panel (E) the power was somewhat increased for an additional measurement (triangles) that emphasizes such a weak line.

Interference of Interacting Matter Waves

Mattias Gustavsson¹, Elmar Haller¹, Manfred J. Mark¹, Johann G. Danzl¹, Russell Hart¹, Andrew J. Daley^{2,3}, Hanns-Christoph Nägerl¹

> ¹Institut für Experimentalphysik und Zentrum für Quantenphysik, Universität Innsbruck, Technikerstraße 25, 6020 Innsbruck, Austria

> ²Institut für Theoretische Physik und Zentrum für Quantenphysik, Universität Innsbruck, Technikerstraße 25, 6020 Innsbruck, Austria

³Institut für Quantenoptik und Quanteninformation der Österreichischen Akademie der Wissenschaften, Technikerstraße 21a, 6020 Innsbruck, Austria

The phenomenon of matter wave interference lies at the heart of quantum physics. It has been observed in various contexts in the limit of non-interacting particles as a single particle effect. Here we observe and control matter wave interference that is driven by interparticle interactions. In our matter wave interferometer, the macroscopic many-body wave function of an interacting atomic Bose-Einstein condensate develops a regular interference pattern, allowing us to detect and directly visualize the effect of interaction-induced phase shifts as time progresses. We demonstrate in a matter wave spin-echo-type experiment that the nonlinear phase evolution is a highly coherent process.

Matter wave interference has been observed since de Broglie's postulate and the develop-

ment of quantum mechanics for various systems, e.g. for electrons (1), neutrons (2), atoms and molecules (3). Macroscopic matter wave interference was first directly observed in the case of two independent atomic Bose-Einstein condensates (BEC) that were brought to overlap (4). This experiment established the notion of the BEC as a macroscopic matter wave and coined the expression of the atom laser in analogy to the laser for the case of photons. Matter wave interferometry (5, 6) exploits matter wave interference to sensitively map relative phase shifts imparted on the de Broglie waves onto population differences, which can then be directly measured. In principle, a crucial role for all matter wave interferometers is played by particle-particle interactions (6, 10). It is commonly expected that such interactions lead to uncontrollable phase shifts and dephasing and form a source of decoherence. Matter wave interferometers, in particular for applications to precision measurements, are thus operated in the dilute single particle limit (7-9). BEC-based atom interferometers (11) are expected to benefit from the extremely low momentum spread and exceptional brightness of the BEC, but they readily enter the nonlinear matter wave regime as a result of the interaction-induced mean field potential. A possible solution is to operate BEC-based interferometers in the non-interacting limit (12, 13) by exploiting the cancellation of the scattering phase shift near a scattering resonance. This condition, however, is difficult or impossible to fulfill for most atomic species.

Here we demonstrate a BEC-based multipath atom interferometer where the dynamics is dominated by interaction-induced phase shifts. We realize the multipath interferometer by loading an interacting BEC into an optical lattice potential along one dimension, coherently splitting the BEC into several parts that are then each subject to different linear and nonlinear phase shifts. The linear phase shifts due to the gravitational force lead to the the well-known phenomenon of Bloch oscillations (*14*, *15*), while the interaction-induced nonlinear phase shifts cause the macroscopic momentum wave function to first spread as a function of time and then, surprisingly, to exhibit high-contrast interference. We demonstrate a high degree of coherence

by reversing the nonlinear phase evolution and refocusing the BEC momentum wave function. A crucial ingredient of our experiments is the capability to tune a, the atomic scattering length which determines the strength of the interaction, by means of a Feshbach resonance (16). In particular, a can be switched to zero to stop the interaction driven part of the evolution in the interferometer or to perform high resolution wave function imaging in momentum space.

When our BEC is loaded into the 1D optical lattice with spacing d, the main features of the system dynamics are largely described, as we will show below in our experiment, by a discrete nonlinear equation (DNLE) in one dimension (17). In brief, this equation can be obtained by expanding the condensate wave function from the Gross-Pitaevskii equation, Ψ , in a basis of wavefunctions $\Psi_j(z, r_{\perp})$ centered at individual lattice sites j, $\Psi(z, r_{\perp}, t) = \sum_j c_j(t) \Psi_j(z, r_{\perp})$. Here, z is the coordinate along the (vertical) lattice direction, r_{\perp} is the transverse coordinate, and $c_j(t)$ are time-dependent complex amplitudes. Due to the small interaction energies in our system, the atoms are restricted to move in the lowest Bloch band and we can write $\Psi_j(r_{\perp}, z) = w_0^{(j)}(z)\Phi_{\perp}(\rho_j, r_{\perp})$, where $w_0^{(j)}(z)$ are the lowest-band Wannier functions localized at the j-th site and $\Phi_{\perp}(\rho_j, r_{\perp})$ is a radial wave function depending on the peak density ρ_j at each site (17). By inserting this form into the Gross-Pitaevskii equation and integrating out the radial direction, the DNLE is obtained,

$$i\hbar \frac{\partial c_j}{\partial t} = J(c_{j-1} + c_{j+1}) + E_j^{\text{int}}(c_j)c_j + V_jc_j.$$

$$\tag{1}$$

Here, J/\hbar is the tunneling rate between neighboring lattice sites, $V_j = Fd j + V_j^{\text{trap}}$ describes the combination of a linear potential with force F and an external, possibly time-varying trapping potential V_j^{trap} , and $E_j^{\text{int}}(c_j)$ is the nonlinear term due to interactions.

In our experiment, we first load the BEC into the vertical lattice and then allow the gravitational force to tilt the lattice potential. We thus enter the limit $Fd \gg J$, in which tunneling between sites is inhibited and the on-site occupation numbers $|c_j|^2$ are constant, determined by the initial density distribution. The time-dependence of the system is then given by the relative phases of each c_j , and the time-dependent 1D wave function $\tilde{\Psi}(q, t)$ in quasi-momentum space q acquires a particularly simple form (18):

$$\widetilde{\Psi}(q,t) = \sum_{j} c_{j}(t) e^{-iqjd} = \sum_{j} c_{j}(0) e^{-i(Fdj + V_{j}^{\text{trap}} + E_{j}^{\text{int}})t/\hbar} e^{-iqjd}$$

$$= \sum_{j} c_{j}(0) e^{-i(q + \frac{Ft}{\hbar})jd} e^{-i(\beta_{\text{tr}}(j-\delta)^{2} - \alpha_{\text{int}}(j-\delta)^{2})t/\hbar}$$
(2)

Here, we have assumed that the external potential is harmonic and is given by $V_j^{\text{trap}} = \beta_{\text{tr}}(j - \delta)^2$, where $\beta_{\text{tr}} = m\omega_{\text{tr}}^2 d^2/2$ characterizes the strength of the potential with trapping frequency ω_{tr} along z for a particle with mass m, and $\delta \in [0, 1]$ describes a possible offset of the potential center with respect to one of the lattice minima along the z-direction. For the interaction term, the spatial dependence is also parabolic, reflecting the fact that we initially load a BEC in the Thomas-Fermi regime. Hence, initially $\alpha_{\text{int}} = m\omega_{\text{lo}}^2 d^2/2$, where ω_{lo} is the trapping frequency during loading of the lattice, as long as the scattering length is kept constant. In the experiments below, δ is the only parameter that we do not fully control. It is constant on the timescale of a single experimental run, but it drifts over the course of minutes as the position of the horizontally propagating laser beam generating the trapping potential is not actively stabilized.

Equation (2) has a simple interpretation. The terms in the exponent linear in j result in Bloch oscillations (14, 15) with angular frequency Fd/\hbar (19). The nonlinear components proportional to j^2 lead to a dephasing of Bloch oscillations, resulting in a time-varying interference pattern for the macroscopic matter wave, as we will demonstrate below. The key here in our experiments is that we have full control over these nonlinear terms, not only over β_{tr} via the external trapping potential, but also over the interaction term characterized by α_{int} , via the initial density distribution, and, more importantly, via the scattering length a. By tuning the scattering length from its initial value a to a' (16), we can ramp α_{int} to a new value α'_{int} , which can in particular be set to zero for a=0 (20). Nonlinear phase terms for matter waves are well known in single particle quantum mechanics. They play an important role for matter wave Talbot interferences (5, 21) and can be visualized in terms of so-called matter wave quantum carpets (22). In these contexts, the phase terms arise from propagation. In our case, the nonlinear phase terms for $\alpha_{int} \neq 0$ arise from interactions and thus lead to a density dependent many-body effect in the multipath atom interferometer.

Our experimental approach initially follows the procedure described in Ref. (12). In brief, within 10 s we produce an essentially pure, tunable BEC (16) in the Thomas-Fermi limit with up to 1.5×10^5 Cs atoms in a crossed-beam dipole trap generated by a vertically (L₁) and a more tightly focused horizontally (L₂) propagating laser beam. The BEC is cigar-shaped with the long axis oriented along the direction of L₂. The trap frequencies are $(\omega_x, \omega_y, \omega_z) = 2\pi \times (39, 5, 39)$ Hz, where x denotes the horizontal direction perpendicular to L_2 , y is the axial direction along L_2 , and z is the vertical direction. We magnetically control a in the range between $0 a_0$ and $300 a_0$ with a resolution of about $0.1 a_0$ (12), where a_0 is the Bohr radius. For BEC production, we work at $a = 210 a_0$. Initially, we support the optical trapping by magnetic levitation against gravity (16). As shown in Figure 1A we superimpose the optical lattice with $d = \lambda/2 = \pi/k$ along the vertical direction, where $\lambda = 1064.5$ nm is the wavelength of the lattice light and k = $2\pi/\lambda$ its wavenumber. To load the BEC into the lattice, we stiffen the horizontal confinement within 1 s, leading to trap frequencies of $2\pi \times (41, 13, 39)$ Hz, and at the same time turn on the lattice potential exponentially to a depth of $8E_R$. Here, $E_R = h^2/(2m\lambda^2) = k_B \times 64 \,\mathrm{nK}$ is the photon recoil energy and m the mass of the Cs atom. The BEC is thus gently loaded into the lattice, occupying about 25 to 35 lattice sites, with up to 7000 atoms at the central site.

We effectively start the multipath atom interferometer and hence the evolution of the interacting macroscopic wave function when we induce Bloch oscillations in the lowest band of the lattice by turning off magnetic levitation and ramping down the power in L₂ within 0.3 ms. With $Fd/\hbar \approx 2\pi \times 1740$ Hz and $J/\hbar \approx 2\pi \times 40$ Hz the on-site occupation numbers $|c_j|^2$ are fixed to their initial values. After an evolution time τ , we close the interferometer by ramping down the lattice in 1 ms and directly image the (vertical) quasi-momentum distribution in the first Brillouin zone (BZ). The ramp is adiabatic with respect to the bandgap and maps quasimomentum onto real momentum (23), which is measured by taking an absorption image on a CCD camera after a period of free expansion. We minimize broadening of the distribution as a result of interactions by switching a to zero during the release and the initial expansion (12). Figure 1B shows absorption images with the characteristic cycling pattern for the first Bloch cycle (14). The Bloch period is about 0.58 ms. The peaks have a root mean square (rms) width of $0.2\hbar k$ and are thus well separated. Figure 1C shows the evolution of the quasi-momentum distribution for cycle phase $\phi = 0$, corresponding to the first image in Figure 1B, for the case of an interacting BEC with $a = 190 a_0$ at an initial peak density of $n = 4 \times 10^{13}$ atoms/cm³, occupying about 35 lattice sites after loading. The wave function spreads out in the BZ within about N = 18 cycles. Surprisingly, after about N = 22 cycles, an interference pattern is observable, which, as time progresses, acquires a maximum contrast of typically 75%. The pattern becomes visible after extended time-of-flight. Figure 1D shows how the contrast emerges for N = 40 cycles. It takes more than 100 ms of expansion for the interference pattern to acquire full contrast (24). Magnetic levitation is needed during the expansion to prevent the BEC being accelerated by gravity and falling out of the field of view. Our imaging technique allows us to resolve structure in momentum space on a scale below $0.1\hbar k$ in a single shot absorption image.

To study this matter wave interference pattern we further follow the evolution of the wave function by taking successive snap shots of the quasi-momentum distribution after completion of an integer number of Bloch cycles. Figure 2A shows how the interference pattern changes in time. It first develops at the edges of the BZ and later becomes clearly visible at the center of the BZ, while the number of interference maxima and minima changes as time progresses. For the chosen parameters, we can follow the evolution of the pattern for more than N = 100 Bloch cycles, corresponding to times beyond 60 ms. This is about a factor 10 longer than the timescale for the initial broadening. We find that the number of maxima and minima and their location in the interference pattern as measured after fixed evolution time τ depend on the initial atomic density, on the strength of the interaction, and on the number of occupied lattice sites. We also find that the measured quasi-momentum distribution for a given τ is reproducible from one experimental realization to the next, except that the pattern appears slightly shifted within the BZ over the course of minutes. We attribute this to a drift of δ , the offset of the lattice minima from the dipole trap center, which leads to a small change of the Bloch frequency (19). We do not actively stabilize the vertical position of L₂ with respect to the lattice, and hence temperature variations in the laboratory slowly change δ .

To understand the interference structure and its evolution in time, we apply the simple analytical model from Equation (2), which holds in the limit J = 0, for our experimental parameters. A numerical integration of the one-dimensional DNLE given in Equation (1) gives essentially identical results and confirms the validity of the J = 0 limit. Figure 2B shows the interference pattern according to the analytical model. The qualitative evolution determined by the model is very well reproduced, when we reduce α_{int} by a factor of 0.9 compared to the value deduced from our experimental parameters. This scale factor accounts primarily for the fact that our simple one-dimensional model does not take into account any horizontal dynamics. In particular, switching off L₂ when starting the evolution leads to excitation of a radial breathing mode in the horizontal plane, reducing the density at each site and modulating it in time. To a first approximation, rescaling of α_{int} accounts for this. Nevertheless, the agreement between our experiment and the analytical model indicates that the dominant driving mechanism for the wave function spreading and interference is the result of the nonlinear phase evolution. In particular, phase coherence is not lost, in contrast to previous experiments (25). We test this coherence and the applicability of our simple analytical model in more detail in two experiments below.

First, Equation (2) suggests that the effect of interactions can be cancelled by application of the initial loading potential. Indeed, cancellation of dephasing by choosing $\alpha_{int} \approx \beta_{tr}$ allows us to observe persistent (linear) Bloch oscillations for an interacting BEC. Figure 3A and 3B show the quasi-momentum distribution as a function of the power in L₂, which determines β_{tr} , after 40 Bloch cycles and cycle phase $\theta = \pi$ and $\theta = 0$, respectively. When the external potential does not balance the effect of interactions ($\alpha_{int} \neq \beta_{tr}$), the condensate wave function is dephased and spread over the whole BZ in less than N = 20 cycles as before. When it agrees with the initial loading potential, Bloch oscillations are clearly visible. Compared with the case without a compensating potential, we now see that the time over which Bloch oscillations can be observed is greatly extended. The transition from a dephased to a non-dephased wave function as a function of confinement strength is quantified in Figure 3C, where the rms-width Δp of the singly-peaked quasi-momentum distribution for N = 40 is plotted as a function of the laser power in L2. Figures 3D and 3E show the time evolution of the quasi-momentum distribution without and with the compensating potential while all other parameters are kept the same. Figure 3D essentially shows the initial broadening of the distribution as described before. Interestingly, the condensate wave function as shown in Figure 3E dephases in a completely different way. Initially, the central peak shows no broadening. It is slowly depopulated, while a much broader background distribution is increasingly populated. After about 100 oscillations, the shape of the central peak starts to develop side lobes or splits in two, with the exact shape varying from one experimental run to the next. The timescale for the loss of interference is a factor 10 larger than the timescale on which the dephasing and hence the initial broadening takes place in the uncompensated case.

Second, we perform a matter wave spin-echo-type experiment. We initially proceed as shown in Figure 2, letting the wave function evolve for a time corresponding to about N = 40Bloch cycles until it is fully dephased and shows, upon measurement, a regular interference structure. We then essentially remove the effect of interactions by ramping to $a = 10 a_0$ within 10 ms. By not switching the interaction entirely off and by ramping comparatively slowly we avoid excessive excitation of the radial breathing mode as a result of the change in the mean field potential at each site. At the same time, we gradually turn on the harmonic potential as given by the horizontal dipole trapping laser beam L₂ within 4 ms to approximately the same depth as during the initial BEC loading phase. From Equation (2) we expect that the wave function now experiences a phase shift with a quadratic spatial dependence with opposite sign, allowing us to reverse the evolution and to recover the initial condition. Figure 4 shows the resulting quasi-momentum distributions. As time progresses, the wave function indeed refocuses while it continues to perform Bloch oscillations. As we do not control the value of δ for a particular run, we record about 10 distributions for each evolution time and select those that are symmetrical, corresponding to cycle phase $\phi = 0$ or $\phi = \pi$. For the chosen strength of the potential, refocusing happens after about 24 Bloch cycles after the ramp of a. This confirms that the initial broadening and dephasing mechanism must have been coherent. We note that we cannot avoid some excitation of the radial breathing mode as seen in the absorption images given in Figure 4.

Our results raise several important questions: To what extent can matter wave interferometry be performed in the presence of interactions? What sets the timescale for the eventual loss of interference contrast? Certainly, our simple analytic model does not predict any loss of contrast. In particular, it should be possible to completely eliminate the effect of interactions with the compensating external potential. However, there are several effects not included in the model that could cause the residual dephasing we observe. Motion in the radial direction, which causes the density and therefore the interaction energy to change over time, could lead to mixing of the different degrees of freedom and hence to additional dephasing. This might apply to our matter wave spin-echo experiment, but in the experiment where we compensate by means of the external potential there is hardly any radial excitation and this effect should not play a role. The appearance of dynamical instabilities (26-28) can be ruled out, as our experimental parameters are outside the unstable region. Going beyond the mean-field treatment, a variety of factors can lead to dephasing. For example, at each lattice site there exists a superposition of number states, accumulating different phases corresponding to their respective interaction energies (29, 30). This leads to an effective dephasing, as the phase on a particular lattice site becomes ill-defined. Basic estimates (29, 30) indicate a dephasing time of about 130 ms for our system, on the same order as we observe.

These experiments constitute a clear demonstration of coherent dynamics in an interacting macroscopic quantum system. This coherence affords a large degree of control over the system, as demonstrated by the possibility to rephase the wave function using an external potential in order to reverse dephasing due to interactions. The control demonstrated here has potential application in matter-wave interferometry, and such a degree of control over the mean-field evolution also opens the possibility to probe beyond-mean-field effects in atom interferometers.

References and Notes

- 1. C. J. Davisson, and L. H. Germer, Nature 119, 558 (1927).
- 2. H. v. Halban, Jr. and P. Preiswerk, C.R. Acad. Sci. 203, 73 (1936).
- 3. I. Estermann and O. Stern, Z. Phys. 61, 95 (1930).
- M. R. Andrews, C. G. Townsend, H.-J. Miesner, D. S. Durfee, D. M. Kurn, W. Ketterle, Science 275, 637 (1997).
- 5. P. Berman (ed.), Atom Interferometry (Academic, 1997).
- 6. A. D. Cronin, J. Schmiedmayer, D. E. Pritchard, arXiv:0712.3703 (2007).
- 7. A. Wicht, J. M. Hensley, E. Sarajlic, S. Chu, Phys. Scr. T102, 82 (2002).
- P. Cladé, E. de Mirandes, M. Cadoret, S. Guellati-Khélifa, C. Schwob, F. Nez, L. Julien, F. Biraben, *Phys. Rev. Lett.* 96, 033001 (2006).
- 9. J. B. Fixler, G. T. Foster, J. M. McGuirk, M. A. Kasevich, Science 315, 74 (2007).
- 10. R. A. Hart, X. Xu, R. Legere, K. Gibble, Nature 446, 892 (2007).
- S. Gupta, K. Dieckmann, Z. Hadzibabic, D. E. Pritchard, *Phys. Rev. Lett.*, **89** 140401 (2002).
- M. Gustavsson, E. Haller, M. J. Mark, J. G. Danzl, G. Rojas-Kopeinig, H.-C. Nägerl, *Phys. Rev. Lett.* 100, 080404 (2008).
- M. Fattori, C. D'Errico, G. Roati, M. Zaccanti, M. Jona-Lasinio, M. Modugno, M. Inguscio,
 G. Modugno, *Phys. Rev. Lett.* 100, 080405 (2008).

- 14. M. Ben Dahan, E. Peik, J. Reichel, Y. Castin, C. Salomon, Phys. Rev. Lett. 76, 4508 (1996).
- 15. B. P. Anderson and M. A. Kasevich, Science 282, 1686 (1998).
- 16. T. Weber, J. Herbig, M. Mark, H.-C. Nägerl, R. Grimm, Science 299, 232 (2003).
- 17. A. Smerzi and A. Trombettoni, Phys. Rev. A 68, 023613 (2003).
- 18. D. Witthaut, M. Werder, S. Mossmann, H. J. Korsch, Phys. Rev. E 71, 036625 (2005).
- 19. This frequency is modified by the presence of the j^2 terms, to $[Fd + 2(\beta_{trap}\delta \alpha_{int}\delta)]/\hbar$.
- 20. The dependence of α'_{int} on a' is not trivial. For the simple case of adiabatic tuning of the scattering length, one has $\alpha'_{\text{int}} = \sqrt{\frac{a'}{a}} \alpha_{\text{int}}$ due to the modification of the two-dimensional radial wave function (17).
- L. Deng, E. W. Hagley, J. Denschlag, J. E. Simsarian, M. Edwards, C. W. Clark, K. Helmerson, S. L. Rolston, W. D. Phillips, *Phys. Rev. Lett.* 83, 5407 (1999).
- 22. A. E. Kaplan, I. Marzoli, W. E. Lamb, Jr, W. P. Schleich, Phys. Rev. A 61, 032101 (2000).
- 23. A. Kastberg, W.D. Phillips, S.L. Rolston, R.J.C. Spreeuw, P.S. Jessen, *Phys. Rev. Lett.* 74, 1542 (1995).
- 24. In general, we find that the contrast is improved when the horizontally confining beam L_1 is not switched off abruptly but is ramped down slowly within the first 55 ms of time-of-flight, reducing the horizontal expansion rate. However, this happens at the cost of some additional momentum broadening along the vertical direction.
- O. Morsch, J. H. Müller, D. Ciampini, M. Cristiani, P. B. Blakie, C. J. Williams, P. S. Julienne, E. Arimondo, *Phys. Rev. A* 67, 031603 (2003).

- 26. Y. Zheng, M. Koštrun, J. Javanainen, Phys. Rev. Lett. 93, 230401 (2004).
- 27. M. Cristiani, O. Morsch, N. Malossi, M. Jona-Lasinio, M. Anderlini, E. Courtade, E. Arimondo, *Opt. Express* **12**, 4 (2004).
- L. Fallani, L. De Sarlo, J. E. Lye, M. Modugno, R. Saers, C. Fort, M. Inguscio, *Phys. Rev. Lett.* 93, 140406 (2004).
- 29. W. Li, A. K. Tuchman, H.-C. Chien, M. A. Kasevich, Phys. Rev. Lett. 98, 040402 (2007).
- 30. A. Imamoglu, M. Lewenstein, L. You, Phys. Rev. Lett. 78, 2511 (1997).
- 31. We thank E. Arimondo, O. Morsch, W. Schleich, A. Smerzi, and A. Buchleitner and his group for helpful discussions. We are indebted to R. Grimm for generous support and gratefully acknowledge funding by the Austrian Ministry of Science and Research (Bundesministerium für Wissenschaft und Forschung) and the Austrian Science Fund (Fonds zur Förderung der wissenschaftlichen Forschung) in form of a START prize grant and through SFB 15. R.H. is supported by a Marie Curie International Incoming Fellowship within the 7th European Community Framework Programme.



Fig. 1. A Experimental configuration: The tunable BEC is formed at the intersection of the vertical guide laser beam L_1 and a horizontal trapping beam L_2 . The lattice is oriented along the vertical direction. Gravity, g, is initially compensated by means of a magnetic force gradient, ∇B . **B** Imaging the first Brillouin zone (BZ): One cycle of Bloch oscillations for a non-interacting BEC as seen in time-of-flight absorption imaging showing narrow peaks cycling through quasi-momentum space for cycle phases $\phi = 0$, $\pi/4$, $\pi/2$, ..., to 2π . **C** Broadening of the distribution and development of an interference pattern for an interacting BEC for $\phi = 0$ after N = 2 to 26 Bloch cycles. **D** Contrast of matter wave interference emerging during time-of-flight expansion for a BEC that completely fills the BZ after N = 40 Bloch cycles. We

define the contrast as $(I_{max} - I_{min})/(I_{max} + I_{min})$, where I_{max} (I_{min}) is the average value of the maxima (minima) of the central peak structure. The error bars are the 1σ statistical error as each data point is the average contrast of 10 experimental runs. The insets show measured quasi-momentum distributions integrated along the transverse direction at two expansion times as indicated.



Fig. 2. Interaction induced macroscopic matter wave interference. A Experimental results showing the quasi-momentum distribution as a function of time given in units of the Bloch period. The absorption images are taken in steps of 4 Bloch cycles for a BEC with an initial peak density of $n=4 \times 10^{13}$ atoms/cm³ loaded into about 35 lattice sites with a=190 a₀. Each image corresponds to a single realization of the experiment. **B** The quasi-momentum wave function evolution according to Equation (2) for $\beta_{tr} = 0$ (no external trap) for $n = 4 \times 10^{13}$ atoms/cm³ loaded into 35 lattice sites with a = 190 a₀. The images are obtained from the 1D quasi-momentum distribution combined with the appropriate site-dependent Thomas-Fermi profiles in the radial direction. α_{int} is slightly rescaled to account for the reduction in density due to transversal dynamics, see text. In A, some additional broadening, largely due to the presence of the horizontal trapping potential during expansion (24), can be seen.



Fig. 3. Cancellation of interaction induced dephasing and observation of persistent Bloch oscillations. Absorption images showing the quasi-momentum distribution for cycle phase $\theta = \pi$ (A) and $\theta = 0$ (B) after N = 40 Bloch cycles and (C) momentum width Δp for $\theta = 0$ as a function of confinement strength, normalized to the confinement strength at loading. D Momentum distribution for $\theta = 0$ as a function of the number N of Bloch cycles when no compensating potential is present, showing fast broadening. E The evolution of the momentum distribution for the case of optimum cancellation of interactions.



Fig. 4. Matter wave spin-echo-type experiment: Rephasing of the BEC from a fully dephased wave function back into a narrow distribution after switching interactions to near zero and turning on an external potential. Time progresses from front to back. The black solid lines correspond to selected quasi-momentum distributions that refocus into the characteristic singly-peaked distribution (cycle phase $\phi = 0$), see text. They are separated in time by 1.15 ms or about two Bloch cycles, and they are offset for clarity. The red solid lines correspond to selected distributions that refocus into the characteristic double-peaked distribution (cycle phase $\phi = \pi$). The images are absorption images corresponding to the adjacent quasi-momentum distributions. Some radial excitation can be seen.