Experiments on two-component quantum gases on an atom chip
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Chapter 5

Controlling one-dimensional spin motion with state-dependent potentials

5.1 Introduction

In this chapter, we present the main results of this thesis. These results have also been submitted for publication [123]. We show that radio-frequency-dressed potentials on atom chips offer a new way to tune the effective interactions in 1D and to control spin motion. We make use of the fact that, for elliptical RF polarizations, different hyperfine states experience different dressed potentials, allowing for state-dependent manipulation [124]. Here we exploit the dependence of the 1D coupling strength on the transverse confinement frequency $\omega_\perp$ [34]. By tuning the transverse confinement for the two states independently through the RF polarization and amplitude, we show that it is possible to control the interactions in a state- and time-dependent manner. Suddenly changing interactions, combined with the state dependence of the axial trapping then results in dynamical evolution in the spin degree of freedom. In particular, we are able to tune to (i) the point where the spin motion is frozen, and (ii) the point where the 1D interactions become spin-independent.

5.2 Experimental procedure

We first discuss our results on the one-dimensional non-equilibrium dynamics for state-independent potentials, highlighting the importance of small differences in interaction parameters. The experimental procedure of the experiments detailed in this chapter is similar to the one used for the experiments of chapter 4. In fact, the individual steps of the two procedures are identical up to an intermediate step in a series of evaporative cooling stages. Whereas in chapter 4 we first mixed the states $|1\rangle$ and $|2\rangle$ when the gas could still be considered non-degenerate and then continued evaporative cooling to degeneracy, here all cooling is performed prior to creating a two-component system. We provide a documented example of the program code used to generate a two-component 1D Bose gas in Appendix A.

The starting point of our experiments is a nearly-pure 1D quasi-condensate in the $|1\rangle = |F = 1, m_f = -1\rangle$ state of $^{87}$Rb in a highly elongated magnetic trap created by an atom chip. To prepare this state, initial evaporative cooling to 120 kHz above the trap bottom is performed (see chapter 4.2), leading to a cold, non-degenerate gas. After this, we apply a short RF pulse with a frequency of 2.22 MHz. This serves to transfer atoms from $|F = 2, m_f = 2\rangle$ to $|2\rangle = |F = 2, m_f = 1\rangle$, as described in chapter 4.2.
Then, we drive a two-photon transition by applying a combined RF (1 MHz)- and MW (6.83 GHz) field. For a pulse time of 0.5 ms, the atomic population is transferred from $|2\rangle$ to $|1\rangle = |F = 1, m_f = -1\rangle$ with nearly 100% transfer efficiency. The single transitions are detuned by 1.26 MHz from the intermediate $|F = 2, m_f = 0\rangle$ level to avoid transfer to this untrapped state. Subsequently, the trap is tightened by ramping up the current in chip wire 5 and the coils responsible for the bias field within 60 ms. This creates a highly elongated Ioffe-Pritchard microtrap with trap frequencies of $\omega/2\pi = 1.9$ kHz and $\omega/2\pi = 26$ Hz.

In the following step, we perform a second and final RF evaporation sweep to ≈ 25 kHz above the trap bottom in order to achieve quantum degeneracy of state $|1\rangle$. This sweep is accompanied by a MW pulse to remove the still populated fully stretched state $|F = 2, m_f = 2\rangle$ by coupling to the untrapped $|F = 1, m_f = 1\rangle$ state. The peak linear atomic density is $n_1 \lesssim 100$ µm$^{-1}$. In this system, both the temperature and chemical potential are small compared to the radial excitation energy ($\mu, k_B T < \hbar \omega$) and the dynamics are restricted to the axial dimension (1D regime).

From this clean starting point, in which only state $|1\rangle$ has a noticeable population, we induce a sudden transition to a coherent superposition of the $|1\rangle$ and $|2\rangle = |F = 2, m_f = 1\rangle$ hyperfine states via a second two-photon pulse, effectively creating a spin-1/2 system. We recorded Rabi oscillations between the two coupled states by varying the two-photon pulse duration and determining the number of atoms per state. An example of the oscillations, recorded under slightly different conditions, is shown in figure 5.1. From a sine fit to this data we deduce a two-photon Rabi frequency of 1.8 kHz.

![Figure 5.1: Rabi oscillations between states $|1\rangle$ and $|2\rangle$ driven by coherent two-photon transition. The number of atoms in states $|1\rangle$ and $|2\rangle$ is indicated red and blue, respectively. Sine fits to the data yield a Rabi frequency of 1.8 kHz.](image)

The two-photon Rabi frequency for the experiments described here is determined as 1.14 kHz, corresponding to a $\pi/2$-pulse duration of 0.22 ms. This is fast compared to the timescale for axial dynamics, but sufficiently slow to prevent radial excitations. Co-
herence times in excess of 1 s have been measured in this setup via Ramsey spectroscopy of dilute thermal clouds.

After performing a $\pi/2$-pulse, which leads to an equal mixture of atoms in states $|1\rangle$ and $|2\rangle$, the resulting non-equilibrium situation is allowed to evolve. The time evolution of the spin distribution is measured by varying the hold time and subsequent sequential state-dependent absorption imaging.

After a short time-of-flight of typically 1 ms, we image the longitudinal distributions, and obtain the linear densities $n_1$ and $n_2$ of the two states along the length of the trap by integrating along the radial direction. Atoms in state $|2\rangle$ are imaged directly using absorption on the $F = 2, F' = 3$ transition with an exposure time of 30 $\mu$s and an optical resolution of 4.2 $\mu$m. Due to an extra repumping step we find a 20% lower detection efficiency for state $|1\rangle$ and a poorer resolution of $\sim 8 \mu$m due to photon recoil, visible in figure 5.8.

### 5.3 Spin motion

In figure 5.2 we present measurements of the evolution of spin polarization $(n_1 - n_2)$ and the total linear density $(n_1 + n_2)$ as a function of hold time. The spin pattern shows clear dynamical evolution [figure 5.2(a)] whereas the total density remains approximately constant with no significant dynamics [figure 5.2(b)]. The spin dynamics can be interpreted as a “focusing” of state $|2\rangle$ in the presence of state $|1\rangle$, resulting in a negative spin polarization $(n_2 > n_1)$ toward the center of the trap.

We find good agreement with the experimental data using the coupled 1D Gross-Pitaevskii equations (1D-GPE) with solutions also shown in figure 5.2(c,d). The 1D-GPE is obtained by integrating the full 3D-GPE over the transverse ground-state wavefunctions (see equations (2.26)), with interaction parameters derived from the intra- and interstate scattering lengths taken from ref. [98]: $a_{11} = 100.4 \cdot a_0$, $a_{22} = 95.44 \cdot a_0$ and $a_{12} = 98.006 \cdot a_0$, where $a_0$ is the Bohr radius. Generalizing for state-dependent harmonic confinement (as will be relevant below) we obtain for the 1D interaction parameters $u_{ij}$ (equations (2.27) with $\omega_{ij} = \omega_{iz}$):

\[
\begin{align*}
    u_{11} &= 2\hbar\omega_{\perp,1}a_{11}, \\
    u_{22} &= 2\hbar\omega_{\perp,2}a_{22}, \\
    u_{12} &= 4\hbar\frac{\omega_{\perp,1}\omega_{\perp,2}}{(\omega_{\perp,1} + \omega_{\perp,2})}a_{12},
\end{align*}
\]  

(5.1)

with $\omega_{\perp,j}$ the transverse trap frequency for state $|j\rangle$. Similarly we use values for the scaled rate constants for inelastic two-body and three-body losses derived (see equations (2.28)) from the 3D values in ref. [98]. The 1D-GPE simulations reproduce the features of the experiment, i.e. absence of dynamics in the total density and the overall structure of the spin dynamics including the time of maximum state separation around $t \approx 75$ ms.
Figure 5.2: One-dimensional spin dynamics and total density after a sudden transfer of internal-state population, for the case of state-independent trapping potentials. Shown are spin polarization \( n_1 - n_2 \), left) and total linear density \( n_1 + n_2 \), right), as a function of axial position and time after the transfer. Top: experiments, bottom: corresponding simulations resulting from integration of two coupled 1D Gross-Pitaevskii equations (GPE). The spin polarization data clearly shows how \( n_2 \) is focused towards the center (blue), while \( n_1 \) moves towards the sides (red); the total density shows little dynamics. Differences between experiment and simulation can be explained by the limited optical resolution of our imaging system and a small tilt of the trap, leading to a slight spatial asymmetry in the experiments.
5.4 Controlling spin motion in a 1D Bose gas

The decay in atom number on a $\gtrsim 100$ ms timescale is dominated by two-body losses in intrastate interactions and between $|2\rangle$ atoms ($\gamma_{12}$ and $\gamma_{22}$) [98].

The rate of spin focusing/defocusing is critically dependent on the precise differences in 1D interaction strengths for the respective internal states, a fact that is readily confirmed by changing these differences in the simulations. The observed general behavior can be understood as follows: in the initial state (an interacting trapped quantum gas in a single internal state in equilibrium) the repulsive interactions balance the external confining potential. Suddenly transferring a fraction of the population to a second internal state with weaker intra- and interstate interactions results in a net contracting force (a confining effective curvature, $c_2 > 0$ in equations (5.3) below) on the population in this second state that dominates the dynamics in the spin polarization. Because the spin-dependent part of the interactions is relatively small, the dynamics in the total density are dominated by the (relatively large) average scattering length which remains nearly constant. Hence the total density shows only weak dynamics; the focusing in $n_2$ is accommodated by “pushing” $n_1$ to the sides (red in figure 5.2).

5.4 Controlling spin motion in a 1D Bose gas

We now describe the state-dependent radio-frequency-dressed potentials that we use to control the spin motion. We consider near-resonant coupling ($\hbar \omega_{rf} \lesssim g_F \mu_B |B|$) of the RF-field with tuneable polarization determined by the relative phase of two independently controlled RF-fields. A cross section of the wire geometry used is shown in figure 5.3(a). The fields originate from direct digital synthesis (DDS) supplied currents in two wires neighboring the Z-shaped trapping wire [125]. With these two fields we can readily control the ellipticity of the total RF-field at the trap position by controlling the relative phase $\phi$ of the RF currents in the two wires. This includes linear (horizontal and vertical) and circular ($\sigma^\pm$) polarizations.

The corresponding dressed-state potential for state $|j\rangle$ (with $j = 1, 2$) has the form $V_j(x, y, z) = (V_0(x, y, z) - \hbar \omega_{rf})^2 + \hbar^2 \Omega_j^2)^{1/2}$ where $V_0(x, y, z)$ is the bare magnetic (harmonic) potential. The state-dependent part of the potential enters through the coupling Rabi frequency $\Omega_j$ [126, 94] (see equation 2.9), which acts to weaken the overall confinement near the trap bottom by an amount given by the dressing parameter $\delta_j$. Taking the second derivative of the potential $V_j$ around the origin yields new trap frequencies,

$$\tilde{\omega}_{\perp,\parallel}^2 = \delta_j \omega_{\perp,\parallel}^2, \text{ where } \delta_j = \Delta / \sqrt{\Omega_j^2 + \Delta^2},$$

with detuning $\Delta = \omega_L - \omega_{rf}$ and Larmor frequency $\omega_L$.

The state-dependent RF potential is characterized using dressed-state RF spectroscopy with a weak additional RF probe [127, 128]. The potential energy at the trap bottom is characterized by the onset of loss as a function of probe frequency which we fit to extract $V_j(0, 0, 0)$. Figure 5.3(b) shows the measured trap bottom as a function of the dressing...
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Figure 5.3: State-dependent potentials. (a) wire geometry used for state-dependent radio-frequency-dressed traps (as in fig. 2.3). The static quadrupole magnetic field and the two RF-fields are indicated by green, orange and purple arrows respectively. The direction of the bias Ioffe field which defines the quantization axis is into the plane of the figure. (b) Trap bottom determined via dressed state RF spectroscopy as a function of \( \phi \). Data points correspond to the measured trap bottom for state \( |1\rangle \) (red) and state \( |2\rangle \) (blue). Solid and dashed curves are fits to the data. The dash-dotted green line indicates the fitted trap Larmor frequency \( \omega_L/2\pi = 2.25 \text{ MHz} \).

Phase \( \phi \), for \( \omega_{rf} = 2\pi \times 2.20 \text{ MHz} \). The trap bottom varies with RF phase between 2250 kHz and 2700 kHz corresponding to a maximum Rabi frequency of \( \Omega_j/2\pi = 450 \text{ kHz} \). For \( \phi = 0.31\pi \) and \( \phi = 1.68\pi \) the potential is maximally state-dependent, corresponding to the pure circular polarizations \( \sigma^- \) and \( \sigma^+ \) respectively (dressing only state \( |1\rangle \) and only state \( |2\rangle \), respectively). The potentials are state-independent for linear polarization at \( \phi = 0, \pi \) (equal dressing of state \( |1\rangle \) and \( |2\rangle \)). The deviation from a simple \( \sin^2(\phi) \) behavior is due to the wire geometry, as the two RF-fields are not quite orthogonal at the trap position.

A fit to the data (solid and dashed curves) taking into account the wire geometry results in an accurate calibration of the key experimental parameters, in particular the Larmor frequency \( \omega_L = 2.25 \text{ MHz} \), RF-field amplitudes \( b_1 = b_2 = 0.53 \text{ G} \) from the two RF wires and the trap-surface distance of 80 \( \mu \text{m} \).

To control the spin motion we turn on the state-dependent dressing directly after preparing the equal superposition of \( |1\rangle \) and \( |2\rangle \). This is done by ramping up the RF currents through chip wires 3 and 6 within 2 ms. The radio-frequency, \( \omega_{rf} = 2\pi \times 2.20 \text{ MHz} \), is chosen slightly below the trap bottom to prevent the formation of a double-well potential. The ramp time is slow compared to the inverse Larmor frequency and the inverse radial trap frequency, but sudden with respect to any axial motion. We use the two circular RF polarizations and various RF amplitudes, corresponding to \( 0.8 < \delta_1 < 1, \delta_2 = 1 \) and \( 0.8 < \delta_2 < 1, \delta_1 = 1 \). For each time step we extract the widths of the axial distributions in both states.

Results for the full range of dressing parameters are depicted in figure 5.6. Figure 5.6(a) shows the calculated interaction strengths taken from equation (5.1) as a function of \( \delta_1 \) and \( \delta_2 \). We have compared the measured widths of the distributions as a function of time.
Figure 5.4: 1D density profiles at hold time $t = 44 \text{ ms}$ for (a) no RF dressing, (b) dressing of state $|1\rangle$ alone ($\phi = 0.31\pi$, $\delta_1 = 0.904$) and (c) dressing state $|2\rangle$ alone ($\phi = 1.68\pi$, $\delta_2 = 0.968$). States $|1\rangle$ and $|2\rangle$ are indicated in red and blue, respectively. The experimental data is represented by dots, gaussian fits by thin dashed lines and the profiles resulting from the corresponding coupled 1D-GP simulation are represented by solid lines.
Figure 5.5: Number of atoms in states $|1\rangle$ (red) and $|2\rangle$ (blue) as function of the hold time in the magnetic trap for three different RF dressing parameters. Subfigure (a) corresponds to no RF dressing, (b) to dressing of state $|1\rangle$ alone ($\phi = 0.31\pi, \delta_1 = 0.904$) and (c) shows the data for dressing state $|2\rangle$ alone ($\phi = 1.68\pi, \delta_2 = 0.968$). The solid lines are the corresponding decay curves obtained from the GP simulations.
Figure 5.6: Overview of the possibilities of the state-dependent potentials, as a function of the dressing parameters (left: varying $\delta_2$, with $\delta_1 = 1$; right: varying $\delta_1$ with $\delta_2 = 1$). (a): 1D interaction strengths, $u_{ij}$ normalised by the bare transverse trap frequency $\omega_{\perp,0}$. (b) Widths of the distribution at $t = 44\,\text{ms}$ and (c) scaled effective curvature $c_j/c_0$ at $t = 0$. Red indicates state $|1\rangle$ (and $u_{11}$) and blue state $|2\rangle$ (and $u_{22}$) and in (a) $u_{12}$ is indicated in green. The widths in (b) are obtained by a fit to the experimentally measured density profiles (dots) and to GPE simulation (shaded regions). The shaded areas in (b) represent the effect of shot-to-shot atom number fluctuations in the experiment.
with solutions of the coupled 1D-GPE. An example of Gaussian fits to the experimental data taken at \( t = 44 \text{ ms} \) including the corresponding 1D density profiles resulting from the 1D-GPE simulations is shown in figure 5.4. The evolution of the number of trapped atoms with hold time is given in figure 5.5. The observed decay of the atom numbers matches the simulation based on the numbers in ref. [98]. The widths and the corresponding simulations for one fixed hold time of 44 ms are shown in figure 5.6(b). The measured widths follow the 1D-GPE simulations closely (taking into account the finite optical resolution), with the biggest uncertainties originating from atom number fluctuations which cause the peak linear density to vary between \( 70 \mu \text{m}^{-1} \) and \( 100 \mu \text{m}^{-1} \) throughout the entire data set (systematic uncertainty shown by shaded regions). The solid vertical line at \( \delta_1 = 0.904 \) indicates the point where the difference in interaction strengths is minimized [figure 5.6(a)] with \( u_{11}, u_{22} \) and \( u_{12} \) differing by less than 0.09\% (30 times reduction in differences when compared to the unmodified interactions). These conditions are of interest for comparing to Bethe Ansatz solutions which require spin-independent interactions [27, 47].

In the two-component (“spin-1/2”) 1D Bose gas, the presence of spin-independent (symmetric) interactions is of particular interest. For all interaction strengths (weak and strong) the dispersion relation of spin waves is quadratic here [104, 27], and the low-energy spin velocity vanishes. As a consequence the usual Luttinger-liquid description [129, 30, 130, 40] cannot be applied. However, it is precisely the point where exact Bethe Ansatz methods can be used [27, 47]. Furthermore, it is the point where buoyancy effects vanish and in the weakly interacting (mean-field) regime it also lies on the border that separates miscible and immiscible regimes of binary superfluids [104], see figure 2.5.

### 5.5 Effective potentials

To explain the data we have to consider both the effect of RF-dressing on the collisional interaction strengths as well as the state-dependent modification to the axial potential. A simple analytical description can be obtained using a Thomas-Fermi description near the cloud center where the cloud shape is an inverted parabola. The combination of the state-dependence of the axial trapping frequency and of the interactions can then be expressed as a net harmonic potential characterised by an effective state-dependent curvature \( c_j \).

We solve for the effective curvatures (figure 5.6c) for \( t \gtrsim 0 \) in our experiments in terms of \( \delta_j \), and find

\[
\frac{c_1}{c_0} = \delta_1 - (1 - \beta)\sqrt{\delta_1} - \frac{a_{12}}{a_{11}} \frac{2\beta \sqrt{\delta_1 \delta_2}}{\sqrt{\delta_1 + \delta_2}} \\
\frac{c_2}{c_0} = \delta_2 - \beta \frac{a_{22}}{a_{11}} \sqrt{\delta_2} - \frac{a_{12}}{a_{11}} \frac{2(1 - \beta) \sqrt{\delta_1 \delta_2}}{\sqrt{\delta_1 + \delta_2}}
\]

(5.3)

Here the first term on the right-hand side reflects the modification to the external axial potential and the second and third terms deal with the modified interactions \( u_{ij} \). The axial curvature of the bare potential is \( c_0 = \omega_\parallel^2 / 2 \) and \( \beta \) corresponds to the fraction of the population transferred to state \( |2\rangle \) (\( \beta \approx 1/2 \) for our experiments).
5.6 Outlook

The dashed line at \( \delta_2 = 0.968 \) in figure [5.6] indicates the point where the difference in interactions is compensated by the state-dependent longitudinal potential and \( c_1 = c_2 \). This point is characterized by small equal curvatures of the effective potentials (including interaction energy) for both states [figure [5.6(c)]], which result in frozen spin dynamics. These conditions are important for applications with on-chip atomic clocks, to minimize inhomogeneous broadening due to mean field shifts. For \( \delta_2 < 0.968 \) the difference in interaction strengths is further enhanced and the time evolution of the spin dynamics becomes inverted, with focusing of state \( |1\rangle \) while state \( |2\rangle \) is pushed outward, as is visible in figure [5.6(b)].

Figure [5.7] shows the full time evolution of the spin polarization for two selected RF-dressing parameters. The selected cases are: dressing of state \( |1\rangle \) alone \((\phi = 0.31\pi, \delta_1 = 0.904)\) [figure [5.7(a)] and dressing state \( |2\rangle \) alone \((\phi = 1.68\pi, \delta_2 = 0.968)\) [figure [5.7(b)]], corresponding to the intersection points in figure [5.6(a)] and (c), respectively. Qualitatively state \( |2\rangle \) focuses faster with RF-dressing applied to state \( |1\rangle \), when compared to the case of state independent potentials [figure [5.2(a)]]. Generally, the simulated density profiles reveal a rich and dynamic nonlinear evolution of the spin polarization, reminiscent of filament propagation in optical systems with competing nonlinearities [131]. This is clearly visible in figure [5.7(c)] for example. This detailed structure depends sensitively on the precise values of the dressing. The development and propagation of this fine structure in the spin polarization is partially observed in the experimental data, but is not fully resolved due to the finite imaging resolution. Convolving the simulated profiles with the point-spread function of our imaging system yields excellent agreement with all of the data. With weak dressing of state \( |2\rangle \) \((\delta_2 = 0.968)\) it is possible to freeze spin dynamics altogether such that the two states maintain their overlap and the widths remain constant (apart from a small in-phase quadrupole oscillation and decay from state \( |2\rangle \)), see figure [5.7(b,d)]. A more quantitative representation of the data, showing excellent agreement between experiment and simulation, is given in figure [5.8], where the widths of the two states are shown for different evolution times. Clearly, the focus point can be identified in figure [5.8(a)] around \( t = 75 \) ms and in figure [5.8(b)] around \( t = 30 \) ms, whereas no focussing is present in figure [5.8(c)].

5.6 Outlook

We have shown that by introducing a small state-dependence to the radial trapping potential using RF-dressing we can precisely tune the 1D interaction parameters in a two-component quantum gas by more than 10%, over an experimentally significant range. In our experiments this modification competes with the state dependence of the axial trapping and provides a new “knob” to control spin motion, leading to tuneable nonlinear behavior.

Our method can be naturally extended in several ways. For instance, control over the interactions without the accompanying state-dependence of the axial trapping can be obtained by using one-dimensional box-shaped potentials [125]. By introducing an additional displacement of the transverse potential in a state-dependent way it is possible to further
Figure 5.7: Spatiotemporal behavior of the spin polarization \((n_1 - n_2)\) following the sudden transfer to the state-dependent potentials. States \(|1\rangle\) and \(|2\rangle\) are indicated red and blue, respectively. (a) shows the evolution with RF parameters corresponding to equal inter-atomic interactions \(\delta_1 = 0.904, \delta_2 = 1\) and (b) equal effective potentials \((\delta_1 = 1, \delta_2 = 0.968)\). (c) and (d) show the results of 1D-GPE simulations corresponding to (a) and (b), respectively.
Figure 5.8: Widths of the atomic distributions as function of hold time, for the parameters indicated by the three vertical lines in figure 5.6. Dots are fits to experimental data, lines are results of 1D GPE solutions. States $|1\rangle$ and $|2\rangle$ are indicated red and blue, respectively. In (a) no RF-dressing is applied $\delta_{1,2} = 1$, as in figure 5.2. (b) and (c) correspond to the data in figure 5.7. (b) shows the evolution with RF parameters corresponding to equal inter-atomic interactions ($\delta_1 = 0.904, \delta_2 = 1$), and (c) equal effective potentials ($\delta_1 = 1, \delta_2 = 0.968$).
reduce $u_{12}$, allowing all three interaction parameters to be tuned independently, something that is not generally possible with a magnetically controlled Feshbach resonance.

The observed spin dynamics depend critically on the precise differences in interaction strengths. For $^{87}$Rb, the three relevant scattering lengths are nearly equal and therefore weak dressing is sufficient to tune the system parameters to the point of symmetric interactions or to where the spin dynamics become frozen. Since the RF parameters can be precisely known, such experiments could also allow precision determination of the scattering length differences. More generally, tuning the system parameters around the point of spin-independent interactions strongly affects the dispersion relation of the spin excitations [104]. In particular this allows the spin velocity to be tuned around zero, providing a new handle for the study of spin waves in one-dimensional atomic gases.

Tunable interactions in two-component quantum gases have important applications in the areas of spin-squeezing and quantum metrology [61, 62], and the ability to control spin motion opens new avenues for future studies of quantum coherence in interacting quantum systems [30, 111, 85, 86, 132]. Our current experiments are performed in the weakly interacting 1D regime and at low temperature, and we find that a description based on two coupled 1D Gross-Pitaevskii equations is sufficient to describe our data. The methods presented here to tune interactions are not limited to this regime, however. In particular, we plan to apply these methods to systems with stronger interactions (e.g., by lowering the 1D density) and with higher temperatures. This will provide experimental tests of (and challenges to) more sophisticated theoretical methods, for both equilibrium and non-equilibrium phenomena. For instance, it will be possible to experimentally explore predictions of the thermodynamic Bethe Ansatz for the two-component Bose gas [47] (as already indicated in chapter 4) and to explore quantum quenches in strongly interacting 1D systems by dynamical control over the spin-dependent interactions. Finally, we expect that the experimental control over spin motion and interactions, as demonstrated here, will benefit the realization of spin-“charge” separation in a Bose gas [27, 29].