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Microwave cavity resonance spectroscopy of ultracold plasmas

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We present an *in situ* study of the electron dynamics of an ultracold ⁸⁵Rb plasma using microwave cavity resonance spectroscopy with submicrosecond temporal resolution and vastly improved sensitivity. We observe a three times faster plasma decay than predicted by a self-similar Gaussian expansion model, which might be attributed to collisional microwave heating and plasma-wall interactions. The demonstrated sensitivity enables nondestructive probing of other (nearly) ionized media of fundamental importance as well, such as cold Rydberg gases and antihydrogen plasmas.

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Ultracold neutral plasmas (UNPs) constitute a subfield in plasma physics which emerged in the late 1990s from experiments on laser cooled xenon atoms in a magneto-optical trap (MOT) [1]. UNPs are created by near-threshold photoionization, resulting in kelvin electron and millikelvin ion temperatures, at densities up to 10^{18} m^{-3} [2]. For ions, it allows us to venture into the strongly coupled regime, in which the potential energy of the particles becomes larger than the kinetic energy. This exciting regime was recently explored by Killian and co-workers by laser cooling ions in an ultracold plasma [3].

Strongly coupled plasmas require a new theoretical approach to describe the underlying physics, as spatial correlations between particles arise and effects like Wigner crystallization occur [4], as has been observed in pure ion plasmas [5] and, on a macroscopic scale, in dusty plasmas [6,7]. It is also expected to occur in high density stellar systems such as white dwarfs and in the crust of neutron stars [8]. However, high density systems are difficult to study because the relevant dynamics occurs at timescales from atto- to femtoseconds. For UNPs, this is pico- to microseconds, due to their much lower density, making them the ideal model systems for the study of strongly coupled systems.

Moreover, UNPs are ideal for studying fundamental plasma heating mechanisms, such as disorder-induced heating, three-body recombination (TBR), and collisional and collisionless microwave absorption [9].

Presently, UNPs are studied using charged particle detection, electronic detection, and optical detection. Charged particle detection is a destructive method in which electrons or ions are accelerated towards a multichannel plate (MCP) detector by applying a static electric field. In combination with a strong rf field, electrons can also be resonantly evaporated [10,11] and guided to an MCP, which has been used to observe plasma oscillations [12], Tonks-Dattner resonances [13], and Rydberg atoms created by TBR [14]. The time resolution, limited by the time-of-flight of the charged particles to the MCP, is typically microseconds. Electronic detection is based on applying a weak rf field instead and measuring the change in amplitude and phase of the signal induced on a nearby electrode resulting from plasma oscillations [15]. Although

this method is not destructive, it is not very sensitive since the induced signal is very weak. Ions of Ca and Sr can also be probed by means of laser-induced fluorescence [16] and absorption imaging [17], respectively, which allows their dynamics to be studied *in situ* at very short timescales (typically nanoseconds to microseconds). However, not all ion species have ground state electric-dipole-allowed transitions accessible to lasers and cameras, limiting their diagnostics mainly to charged particle detection [18].

In this Rapid Communication, we present *microwave cavity resonance spectroscopy* (MCRS) as a non-destructive diagnostic to reveal the electron dynamics of ultracold plasmas *in situ*. This technique is based on measuring the frequency shift of a resonant eigenmode inside a microwave cavity, induced by the presence of a UNP. MCRS has been used previously to measure the dynamics of dusty plasmas [19], rf plasmas [20], and plasmas induced by extreme ultraviolet irradiation [21–23]. Here, we demonstrate the creation of an ultracold plasma from a laser cooled ⁸⁵Rb gas inside a microwave cavity, and perform MCRS to determine the electron density of ultracold plasmas as a function of time, without a destructive static electric field. We show that we can detect as few as 3×10^4 electrons, two orders of magnitude more sensitive than earlier MCRS measurements [23], with nanosecond temporal resolution.

MCRS is based on the shift $\Delta\omega$ of the resonance frequency ω_0 of a microwave cavity due to the perturbation of the resonant electromagnetic eigenmode by a dielectric medium inside the cavity, in this case a cloud of plasma. For a nonmagnetized collisionless plasma, the relative permittivity ϵ_r is given by the Drude-Lorentz model: $\epsilon_r = 1 - \omega_{p,e}^2/\omega_0^2$ with $\omega_{p,e} = \sqrt{n_e e^2 / \epsilon_0 m_e}$ the electron plasma frequency, n_e the electron density, e the elementary charge, ϵ_0 the vacuum permittivity, and m_e the electron mass. Using the *Slater perturbation theorem* [24] it follows that

$$\frac{\Delta\omega}{\omega_0} = \frac{e^2}{2\omega_0^2 \epsilon_0 m_e} \bar{n}_e \quad \text{with} \quad \bar{n}_e \equiv \frac{\int_{V_{\text{cav}}} n_e(\mathbf{r}) E^2(\mathbf{r}) d^3\mathbf{r}}{\int_{V_{\text{cav}}} E^2(\mathbf{r}) d^3\mathbf{r}} \quad (1)$$

the field-averaged electron density, $E(\mathbf{r})$ the electric field amplitude at position $\mathbf{r}$, and V_{cav} the cavity volume. If we

FIG. 1. Experimental setup showing the microwave cavity with diffraction grating chip.

assume a Gaussian-shaped ultracold plasma, with rms size σ much smaller than the inner dimensions of the cavity, and a cylindrical cavity with a TM_{010} mode, then

$$\bar{n}_e \approx \frac{N_e}{V_{\text{cav}} J_1^2(x_{01})}. \quad (2)$$

Here, J_n is the n th order Bessel function of the first kind ($n \in \mathbb{N}_0$), x_{01} the first root of J_0 , and N_e is the total number of electrons in the plasma. For a cavity with $\omega_0/2\pi = 5$ GHz, and a height of 1 cm, containing $N_e = 5 \times 10^6$ electrons, this results in $\bar{n}_e = 1 \times 10^{12} \text{ m}^{-3}$ or $\Delta\omega/2\pi = 9$ kHz. This is of the same order of magnitude as the lower detection limit in previous MCRS experiments, which were noise limited [23]. In this work we reduced the noise level by two orders of magnitude and could therefore measure the frequency shift induced by a decaying ultracold plasma.

To create a UNP, we implemented a compact magneto-optical trap based on a diffraction grating chip to laser cool and trap ^{85}Rb atoms (see Fig. 1) [25,26].

We used a $20 \times 20 \text{ mm}^2$ diffraction grating chip with a 100 nm Al coating and an etch depth of 195 nm [27]. The chip has a diffraction angle $\theta = 40.5^\circ$ and a period $d = 1.2 \mu\text{m}$, optimized for the laser cooling wavelength $\lambda = 780 \text{ nm}$ for rubidium [28]. The laser cooling beam has a nearly flat-top profile with a radius $r_b = 12.7 \text{ mm}$ and an intensity $I_0 = 24 \text{ mW/cm}^2$. This creates an overlap volume with a height $h = r_b \cot(\theta) \sim 1 \text{ cm}$, which fits in the 5 GHz cavity and allows us to trap 7×10^7 atoms [25,26].

The cold atomic cloud has a Gaussian density profile with rms sizes $s_x = 1.2 \text{ mm}$, $s_y = 0.8 \text{ mm}$, and $s_z = 0.5 \text{ mm}$. The ultracold plasma is created in a two-step photoionization process, using the 780-nm laser cooling beam to excite the ^{85}Rb atoms from the $5^2S_{1/2}(F=3)$ to the $5^2P_{3/2}(F'=4)$ state, and subsequently, using a 480-nm pulsed dye laser (6 ns) to photoionize the atoms near threshold. The amount of excess energy determines the initial plasma electron temperature T_e . In our case, the ionization laser wavelength was set to 477.86 nm, such that $T_e \approx 50 \text{ K}$.

The dye laser has a $1/e^2$ Gaussian beam diameter of $1.4 \times 0.8 \text{ mm}^2$ and a pulse energy of $260 \mu\text{J}$, such that a large fraction (15%) of the excited atoms in the MOT is ionized. The experiment is repeated at 10 Hz.

Based on the estimation for the field-averaged electron density in Eq. (2), a 5-GHz aluminum microwave cavity was designed for the TM_{010} eigenmode and simulated with CST MICROWAVE STUDIO [29]. The cavity (see Fig. 1) has four small access holes at the sides for access of the ionization laser as well as for determining the atom number and temperature of the MOT by fluorescence imaging measurements. The diffraction grating chip is positioned in a recess in the bottom of the cavity and fixed with two conducting clamps to prevent charge buildup of the plasma on the grating. At the top the cavity has a large opening for the laser cooling beam. In order to reduce power losses, a 2.6-cm tube was added on the top of the opening, which acts a reflecting waveguide for 5 GHz. This results in a high cavity quality factor $Q = 4 \times 10^3$.

However, $Q = 5 \times 10^2$ in the presence of the diffraction grating chip, because the chip only has a 100-nm conducting layer of Al on top of an insulating Si substrate, which increases the cavity's wall resistance. This low Q results in a response time $\tau = Q/\omega_0 = 8 \text{ ns}$ of the cavity, which is much better than the temporal resolution of other electron diagnostics currently used for UNPs.

The rf signal for the cavity's electromagnetic field is produced by a 10-dBm microwave generator, which was set to a frequency $\omega/2\pi$ between 4981 and 4991 MHz. This generator transfers the rf power through a directional coupler to the cavity. A linear antenna in the cavity, next to the diffraction grating chip, couples the rf power electrically to the cavity and picks up the reflected power. A fraction of the reflected power (-10 dB) is split off by the directional coupler and guided to a logarithmic detector. This detector converts the rf signal into a reflection voltage that is subsequently measured on an oscilloscope.

Ideally, the ratio of the absorbed power P_{abs} to the total power P_{tot} transferred to the cavity is described by a Lorentzian function L , i.e., $P_{\text{abs}}/P_{\text{tot}} = L(\omega, \omega_0, I, w)$ with I the height and w the full width-at half maximum of the Lorentzian. Because we used a logarithmic detector, and measured in reflection mode instead of absorption mode, this means that the reflected voltage V_r is described by

$$V_r = V_0 + C \log[1 - L(\omega, \omega_0, I, w)], \quad (3)$$

with V_0 and C constants, depending on the detector configuration. In Fig. 2, the measured reflection voltage is shown as a function of the frequency. By curve fitting with Eq. (3), the resonance frequency of the cavity in vacuum was determined to be $\omega_0/2\pi = 4985.8 \text{ MHz}$.

Subsequently, an ultracold plasma was created by photoionizing the laser cooled ^{85}Rb atoms in the cavity. The reflection voltage was measured and averaged over 100 ionization pulses, as shown in Fig. 3 for nine microwave frequencies. At $t = t_0 = 10 \mu\text{s}$, an ionization pulse creates a plasma, leading to a change in reflection voltage. For frequencies $\omega < \omega_0 + \frac{1}{2}\Delta\omega$ ($\omega > \omega_0 + \frac{1}{2}\Delta\omega$) this change is negative (positive), as the resonance curve V_r in Fig. 2 moves to the right for a positive frequency shift [Eq. (1)].

FIG. 2. Blue curve: measured microwave cavity reflection voltage V_r as a function of applied frequency ω , with curve fit according to Eq. (3). Red curve: measured and calculated change in reflection voltage ΔV_r right after the moment the ionization pulse creates the plasma.

When the reflection voltages of curves with a different frequency are combined for each time t , we can obtain a set of resonance curves, similar to that of an empty cavity (Fig. 2). By fitting each of those curves with Eq. (3), the resonance frequency shift, and thus the field-averaged electron density as a function of time can be determined, as shown in Fig. 4.

At the moment of ionization the resonance frequency shifts by 5.5 kHz, corresponding to a field-averaged density $\bar{n}_e = 7 \times 10^{11} \text{ m}^{-3}$, in agreement with the estimation of Eq. (2). Within $\sim 500 \text{ ns}$, the density decays to $\bar{n}_e = 6 \times 10^{11} \text{ m}^{-3}$, after which the decay slows down substantially. The initial peak, which clearly demonstrates the high temporal resolution of the method (see inset in Fig. 4), is attributed to the escape of fast electrons, which stops when the remaining plasma gets sufficiently positively charged.

The 1σ noise level is 32 Hz, corresponding to a lower detection limit of $\sim 10^9 \text{ m}^{-3}$ (or 3×10^4 electrons), two orders of magnitude less than previous MCRS measurements [23]. This was accomplished by electronic shielding, filtering, averaging, and optimizing the oscilloscope voltage range for each microwave frequency to prevent discretization errors.

Figure 2 shows the measured change in reflection voltage ΔV_r at the moment of photoionization for different frequencies, obtained from the reflection curves in Fig. 3. The solid curve is the calculated change in reflection voltage $\Delta V_r = -(dV_r/d\omega)\Delta\omega$ which agrees well, providing a consistency check of the above analysis. In addition, note that the signal-to-noise ratio is proportional to the maximum derivative $(dV_r/d\omega)_{\text{max}}$ and therefore to Q .

To model the observed decay of the plasma, a self-similar Gaussian expansion model, as described in Ref. [12], was used. The electron density n_e at time $t = t_0$, created by the photoionization of the laser cooled rubidium gas is given by

$$n_e(\mathbf{r}, t = t_0) = f_{\text{exc}} \left[1 - \exp \left(-\frac{\sigma_{\text{PI}} F(\mathbf{r})}{E_{\text{ph}}} \right) \right] n_a(\mathbf{r}), \quad (4)$$

with n_a the ^{85}Rb atomic density, f_{exc} the atomic excited state fraction, $\sigma_{\text{PI}} = 17 \times 10^{-22} \text{ m}^2$ the rubidium photoionization

FIG. 3. Top: microwave cavity reflection voltage V_r as a function of time t for different applied frequencies to the cavity. Bottom: photodiode signal V_{pd} , showing that at $t = t_0 = 10 \mu\text{s}$ the ionization pulse enters the cavity and creates the plasma.

cross section from the $5^2P_{3/2}$ excited state [30], and F and E_{ph} the ionization laser fluence and photon energy, respectively.

The atomic excited state density $f_{\text{exc}} n_a(\mathbf{r})$ was obtained by fluorescence measurements and the fluence of the ionization laser was monitored during measurements using a beam splitter cube.

For small arguments of the exponential in Eq. (4), the electron density is approximately Gaussian, because both the fluence $F(\mathbf{r})$ and atomic density $n_a(\mathbf{r})$ are Gaussian functions. Therefore we fitted the calculated density profile obtained from Eq. (4) with a Gaussian function. This results in a plasma with an initial rms size $\sigma_{0x} = 405 \mu\text{m}$, $\sigma_{0y} = 785 \mu\text{m}$, and $\sigma_{0z} = 218 \mu\text{m}$, with a total of 5.2×10^6 electrons. These numbers correspond to an initial peak density of $4.8 \times 10^{15} \text{ m}^{-3}$.

Immediately after photoionization ($t > t_0$), the self-similar Gaussian plasma model states that the rms sizes $\sigma_i(t)$ of the electron density in directions $i \in \{x, y, z\}$ change as $\sigma_i(t) = \sqrt{\sigma_{0i}^2 + v_{\text{exp}}^2(t - t_0)^2}$. Here v_{exp} is the expansion velocity of the plasma, determined by the excess energy E_{exc} of the ionization laser, i.e., $E_{\text{exc}} \approx \frac{3}{2} m_i v_{\text{exp}}^2$, with m_i the mass of the rubidium ions.

When this model for the electron density is applied together with the electric field map of the cavity obtained from CST MICROWAVE STUDIO simulations, the field-averaged electron density can be calculated with Eq. (1), which is plotted

FIG. 4. Field-averaged electron density as a function of time.

in Fig. 4. It shows that the initial height of the calculated field-averaged electron density matches the measured density quite well.

However, the field-averaged density falls off more quickly in the measurement than what would be expected based on the self-similar Gaussian expansion model; for 50 K electrons, the calculated expansion velocity is 66 m/s. In order to match the plasma expansion time of the model better with the measurement, roughly a factor of 3 higher expansion velocity is required. This might be attributed to heating of the plasma.

Several heating mechanisms have been proposed in the past to explain higher UNP expansion velocities than expected on the basis of the excess energy of the ionization laser. These mechanisms include three-body recombination, continuum lowering, and disorder-induced heating [31,32]. However, in our case, none of these heating mechanisms are sufficient to explain the discrepancy.

One effect that could possibly explain a higher expansion velocity is collisional absorption of microwave energy due to deflection of the quivering electrons in the stationary Coulomb fields of the ions [9]. The amount of energy absorbed by the electrons depends on the electric field amplitude. For a 5 GHz TM_{010} cavity, with an input power $P = 5.7$ dBm, the electric field amplitude $E_0 = \sqrt{2PQ/[\omega_0\epsilon_0 V_{\text{cav}} J_1^2(x_{01})]} = 1.9$ kV/m. The associated quiver velocity of the electrons is $v_{\text{osc}} = eE_0/(m_e\omega_0) = 1.1 \times 10^4$ m/s, which is comparable to their thermal velocity $v_{\text{th}} = \sqrt{k_B T_e/m_e} = 2.8 \times 10^4$ m/s, with k_B the Boltzmann constant. The energy absorbed per unit time by the electrons is $P_{ei} = 2\nu_{ei}U_p$, with $U_p = \frac{1}{4}m_e v_{\text{osc}}^2$ the ponderomotive energy of the electrons in the electric field and ν_{ei} the effective electron-ion collision frequency. For the given field strength in our cavity $\nu_{ei} \approx \nu_s = \sqrt{2/(3\pi)}\omega_{p,e}\Gamma_e^{3/2} \ln \Lambda$, as has been demonstrated by charged particle tracking simulations in Ref. [9], with ν_s the Spitzer collision frequency, Γ_e the electron coupling parameter, and $\ln \Lambda$ the Coulomb logarithm. Calculating the temperature increase per unit of time shows that $P_{ei}/k_B \sim 0.9 \times 10^2$ K/ μs , indicating that the electric field could indeed have a significant heating effect on our ultracold plasma for early times. For later times, the temperature is higher as a consequence of the initial heating and the electron density decreases, causing the effect of collisional heating to become automatically less significant.

Additionally, the presence of the cavity walls could play a role. The walls might lead to deviations in the Gaussian density distribution of the UNP, which causes a breakdown of the self-similarity assumption. For a plasma with an initial temperature of 50 K, this could be important, because its $2\sigma_z$ width reaches the nearest wall (where the diffraction chip is located) in $\sim 30 \mu\text{s}$, roughly half of the lifetime of the field-averaged electron density. For higher expansion velocities, the effect could be even larger.

To conclude, we successfully introduced microwave cavity resonance spectroscopy as a diagnostic enabling *in situ* studies of the electron dynamics of ultracold plasmas, without the need for a destructive static electric field. We created the ultracold plasma from a laser cooled gas of ^{85}Rb atoms inside a compact magneto-optical trap that is based on a diffraction grating chip, placed inside a microwave cavity. We measured the field-averaged electron density as function of time, resulting in an initial field-averaged density of $7 \times 10^{11} \text{ m}^{-3}$, with a temporal resolution ~ 100 ns and a lower detection limit of $\sim 10^9 \text{ m}^{-3}$ (3×10^4 electrons), two orders of magnitude better than earlier MCRS measurements. The initial field-averaged density matches well with the calculated value based on the expected number of electrons and the electric fields in the cavity. However, we observed a faster than expected decay in the density, possibly due to collisional absorption or cavity wall effects.

For future experiments, the MCRS technique could be improved even further by (1) increasing the quality factor. The signal-to-noise ratio is proportional to Q and can therefore be improved by at least an order of magnitude, implying a detection limit of ~ 1 Hz; (2) monitoring frequency shifts of multiple TM eigenmodes instead of only the TM_{010} mode. This allows the evolution of the spatial distribution to be followed in real time, which has been successfully demonstrated in highly transient EUV-induced plasmas [23].

Those improvements will enable the study of the electron dynamics of UNPs *in situ*, independent of atomic species, with unprecedented sensitivity, temporal resolution, and spatially resolved. This could be interesting, in particular, for unraveling competing fundamental plasma heating mechanisms.

In addition, it is intriguing to note that MCRS could also be used to study cold Rydberg gases, owing to their high polarizability. For example, a ^{85}Rb atom in the $43s$ state in a 5 GHz cavity has a dynamic polarizability $\alpha \approx -1.77 \text{ kHz}/(\text{V m}^{-1})^2$ [33]. A basic estimation yields an induced frequency shift per atom $\Delta f \approx -\alpha\omega_0/(\pi\epsilon_0 V_{\text{cav}})$, only one order of magnitude lower than the frequency shift per electron-ion pair in the ultracold plasma.

Finally, MCRS could be an interesting nondestructive diagnostic for antihydrogen ($\bar{\text{H}}$) experiments, allowing the detection of initial high density $\bar{\text{H}}$ plasmas and $\bar{\text{H}}$ Rydberg gases [34,35].

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