THE JAYNES-CUMMINGS MODEL AND THE ONE-ATOM-MASER

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ABSTRACT. In this paper experiments performed with the one-atom maser are reviewed. Furthermore, possible experiments to test basic quantum physics are discussed.

1. Introduction, the One-Atom-Maser

The most promising avenue to study the generation process of radiation in lasers and masers is to drive a maser consisting of a single mode cavity by single atoms. This system, at first glance, seems to be another example of a Gedanken-experiment treated in the pioneering work of Jaynes and Cummings (1963), but such a one-atom maser (Meschede, Walther and Müller, 1985) really exists and can in addition be used to study the basic principles of radiation-atom interaction. The advantages of the system are:

(1) it is the first maser which sustains oscillations with less than one atom on the average in the cavity,
(2) this setup allows to study in detail the conditions necessary to obtain nonclassical radiation, especially radiation with sub-Poissonian photon statistics in a maser system directly in a Poissonian pumping process, and
(3) it is possible to study a variety of phenomena of a quantum field including the quantum measurement process.

What are the tools that make this device work: It was the enormous progress in constructing superconducting cavities together with the laser preparation of highly excited atoms – Rydberg atoms – that have made the realization of such a one-atom maser possible. Rydberg atoms have quite remarkable properties (Haroche and Raimond, 1985; Gallas, Leuchs, Walther and Figger, 1985) which make them ideal for such experiments: The probability of induced transitions between neighboring states of a Rydberg atom scales as n^4, where n denotes the principle quantum number. Consequently, a few photons are enough to saturate the transition between adjacent levels. Moreover, the spontaneous lifetime of a highly excited state is very large. We obtain a maser by injecting these Rydberg atoms into a superconducting cavity with a high quality factor. The injection rate is such that on the average there is less than one atom present inside the resonator at any time. A transition between two neighboring Rydberg levels is resonantly coupled to a single mode of the cavity field. Due to the high quality factor of the cavity, the radiation decay time is much larger than the characteristic time of the atom-field interaction, which is given by the inverse of the single-photon Rabi-frequency. Therefore it is possible to observe the dynamics (Jaynes and Cummings, 1963) of the energy exchange between atom and field mode leading to collapse and revivals in the Rabi oscillations (Eberly, Narozhny and Sanchez-Mondragon, 1980; Rempe, Walther and Klein, 1987). Moreover a field is built up inside the cavity when the mean time between the atoms injected into the cavity is shorter than the cavity decay time.

The detailed experimental setup of the one-atom maser is shown in Figure 1. A highly collimated beam of rubidium atoms passes through a Fizeau velocity selector. Before entering the superconducting cavity, the atoms are excited into the upper maser level $63p_{3/2}$ by the frequency-doubled light of a cw ring dye laser. The laser frequency is stabilized onto the atomic transition $5s_{1/2} \rightarrow 63p_{3/2}$, which has a width determined by the laser linewidth and the transit time broadening corresponding to a total of a few MHz. In this way, it is possible to prepare a very stable beam of excited atoms. The ultraviolet light is linearly polarized parallel to the electric field of the cavity. Therefore only $\Delta m = 0$ transitions are excited by both the laser beam and the microwave field. The superconducting niobium maser cavity is cooled down to a temperature of 0.5 K by means of a ^{3}He cryostat. At such a low temperature the number of thermal photons is reduced to about 0.15 at a frequency of 21.5 GHz. The cryostat is carefully designed to prevent room temperature microwave photons from leaking into the cavity. This would considerably increase the temperature of the radiation field above the temperature of the cavity walls. The quality factor of the cavity is 3×10^{10} corresponding to a photon storage time of about 0.2 s. The cavity is carefully shielded against magnetic fields by several layers of cryoperm. In addition, three pairs of Helmholtz coils are used to compensate the earth's magnetic field to a value of several mG in a volume of $10\times4\times4$ cm^3. This is necessary in order to achieve the high quality factor and prevent the different magnetic substates of the maser levels from mixing during the atom-field interaction time. Two maser transitions from the $63p_{3/2}$ level to the $61d_{3/2}$ and to the $61d_{5/2}$ level are studied.

The Rydberg atoms in the upper and lower maser levels are detected in two separate field ionization detectors. The field strength is adjusted so as to ensure that in the first detector the atoms in the upper level are ionized, but not those in the lower level, they are then ionized in the second field.

To demonstrate maser operation, the cavity is tuned over the $63p_{3/2} - 61d_{3/2}$ transition and the flux of atoms in the excited state is recorded simultaneously. Transitions from the initially prepared $63p_{3/2}$ state to the $61d_{3/2}$ level (21.50658 GHz) are detected by a reduction of the electron count rate.

In the case of measurements at a cavity temperature of 0.5 K, shown in Figure 2, a reduction of the $63p_{3/2}$ signal can be clearly seen for atomic fluxes as small as 1750 atoms/s. An increase in flux causes power broadening and a small shift. This shift is attributed to the ac Stark effect, caused predominantly by virtual transitions to neighboring Rydberg levels. Over the range from 1750 to 28000 atoms/s the field ionization signal at resonance is independent of the particle flux which indicates that the transition is saturated. This, and the observed power broadening show that there is a multiple exchange of photons between Rydberg atoms and the cavity field.

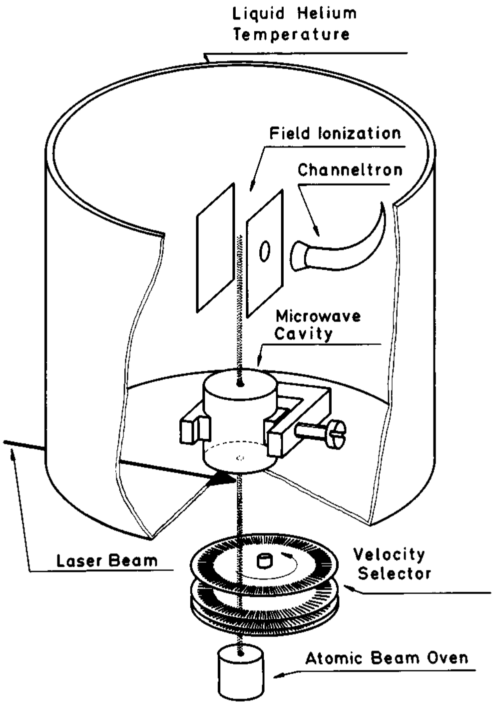


Fig. 1. Scheme of the one-atom maser. To suppress blackbody induced transitions to neighboring states, the Rydberg atoms are excited inside the liquid-Helium-cooled environment.

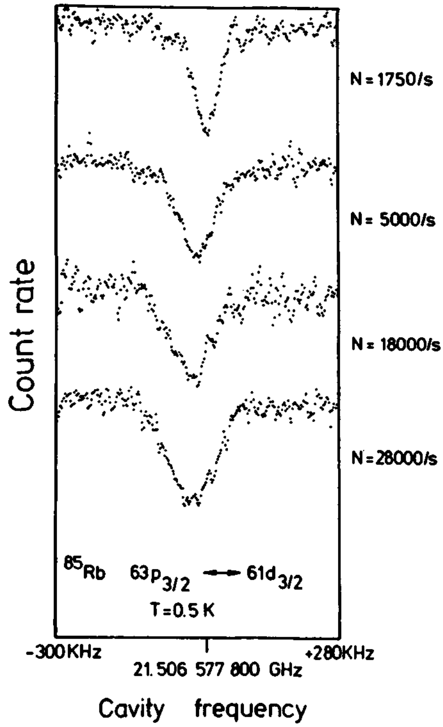
For an average transit time of the Rydberg atoms through the cavity of 50 μs and a flux of 1750 atoms/s we obtain that approximately 0.09 Rydberg atoms are in the cavity on the average. According to Poisson statistics this implies that more than 90% of the events are due to single atoms. This clearly demonstrates that single atoms are able to maintain a continuous oscillation of the cavity with a field corresponding to a mean number of photons between unity and several hundreds.

2. Single Atom Inside a Resonant Cavity - Oscillatory Regime

The experimental set-up described in section 1 is suitable to test the Jaynes-Cummings model. An important requirement is that the atoms of the beam have a homogeneous velocity so that it is possible to observe the Rabi-nutation induced by the cavity field directly. This is not possible with the broad Maxwellian velocity distribution. The Fizeau-type velocity selector is therefore used, so that a fixed atom-field interaction time is obtained (Rempe, Walther, Klein, 1987). Changing the selected velocity leads to a different interaction time and leaves the atom in another phase of the Rabi cycle when it reaches the detector.

The experimental results obtained for the $63p_{3/2} - 61d_{5/2}$ transition are shown in Figure 3. In the figure the ratio between the field ionization signals on and off resonance are plotted versus the interaction time of the atoms in the cavity. The solid curve was calculated using

Fig. 2. A maser transition of the one-atom maser manifests itself in a decrease of atoms in the excited state. The flux of excited atoms N governs the pump intensity. Power broadening of the resonance line demonstrates the multiple exchange of a photon between the cavity field and the atom passing through the resonator. For these experiments no velocity selection of the atoms has been used.



the Jaynes-Cummings model (Jaynes, Cummings, 1963), which is in very good agreement with the experiment. The total uncertainty in the velocity of the atoms is 4% in this measurement. The error in the signal follows from the statistics of the ionization signal and amounts to 4%. The measurement is made with the cavity at 2.5K and $Q = 2.7 \cdot 10^8$ ($\tau_d = 2ms$). There are on the average 2 thermal photons in the cavity. The number of maser photons is small compared with the number of blackbody photons.

The experimental result shown in the upper part of Figure 3 is obtained with very low atom beam flux ($N = 3000s^{-1}$ and $n_m = 0.5$; n_m is the number of photons accumulated in the cavity). When the atomic beam flux is increased more photons are piled up in the cavity. Measurements with $N = 2000s^{-1}$ ($n_m = 2$) and $N = 3000s^{-1}$ ($n_m = 3$) are shown in the lower part of Figure 3. The maximum of $P_e(t)$ at $70\mu s$ flattens with increasing photon number, thus demonstrating the collapse of the Rabi nutation induced by the resonant maser field. Figure 3 (bottom) shows that for atom-field interaction times between $50\mu s$ and about $130\mu s$ $P_e(t)$ does not change as a function of time. Nevertheless, at about $150\mu s$, $P_e(t)$ starts to oscillate again, thus showing the revival predicted by the Jaynes-Cummings model. The variation of the Rabi nutation dynamics with increasing atomic beam fluxes and thus with increasing photon numbers in the cavity generated by stimulated emission is obvious from Figures 3.

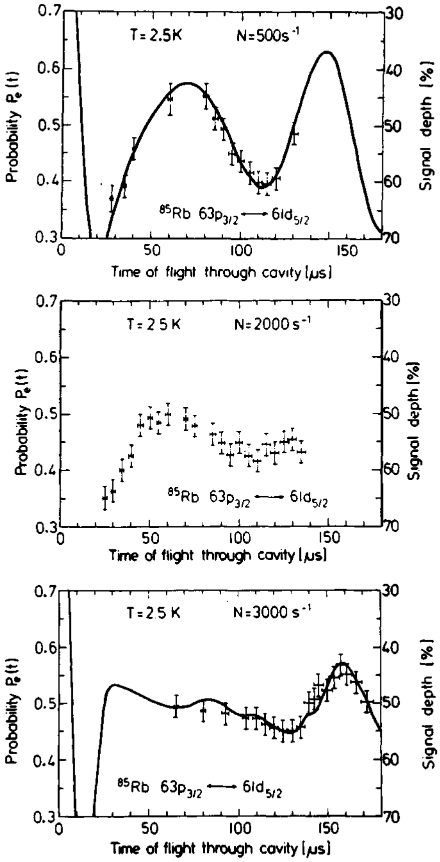


Fig. 3. Measured probability of finding the Rydberg atom in the upper maser level after it passed through the cavity. The cavity is tuned to the $63p_{3/2} - 61d_{5/2}$ transition for increasing flux of the atoms $N = 500s^{-1}$, $N = 2000s^{-1}$ and $N = 3000s^{-1}$. The solid lines represent the theoretical predictions of the Jaynes-Cummings model.

3. A New Source of Nonclassical Light

One of the most interesting questions in connection with the one- atom maser is the photon statistics of the electromagnetic field generated in the superconducting cavity. This problem will be discussed in this section.

Electromagnetic radiation can show nonclassical properties (Walls, 1979, 1983, 1986), that is properties that cannot be explained by classical probability theory. Loosely speaking we need to invoke "negative probabilities" to get deeper insight into these features. We know of essentially three phenomena which demonstrate the nonclassical character of light: photon antibunching (Kimble, Dagenais and Mandel, 1977; Cresser, Häger, Leuchs, Rateike and Walther, 1982), sub-Poissonian photon statistics (Short and Mandel, 1983) and squeezing (Slusher, Hollberg, Yurke, Mertz and Valley, 1985; Loudon and Knight, 1987). Mostly methods of nonlinear optics are employed to generate nonclassical radiation.

However, also the fluorescence light from a single atom caught in a trap exhibits nonclassical features (Carmichael and Walls, 1976; Diedrich and Walther, 1987).

Another nonclassical light generator is the one-atom maser. We recall that the Fizeau velocity selector preselects the velocity of the atoms: Hence the interaction time is well-defined which leads to conditions usually not achievable in standard masers (Filipowicz, Javanainen and Meystre, 1986; Lugiato, Scully and Walther, 1987; Krause, Scully and Walther, 1987; Krause, Scully, Walther and Walther, 1989; Meystre, 1987; Meystre, Rempe and Walther, 1988; Slosser, Meystre and Wright, 1990). This has a very important consequence when the intensity of the maser field grows as more and more atoms give their excitation energy to the field: Even in the absence of dissipation this increase in photon number is stopped when the increasing Rabi-frequency leads to a situation where the atoms reabsorb the photon and leave the cavity in the upper state. For any photon number, this can be achieved by appropriately adjusting the velocity of the atoms. In this case the maser field is not changed any more and the number distribution of the photons in the cavity is sub-Poissonian (Filipowicz, Javanainen and Meystre, 1986; Lugiato, Scully and Walther, 1987), that is narrower than a Poisson distribution. Even a number state that is a state of well-defined photon number can be generated (Krause, Scully and Walther, 1987; Krause, Scully, Walther and Walther, 1989; Meystre, 1987) using a cavity with a high enough quality factor. If there are no thermal photons in the cavity - a condition achievable by cooling the resonator to an extremely low temperature - very interesting features such as trapping states show up (Meystre, Rempe and Walther, 1988). In addition, steady-state macroscopic quantum superpositions can be generated in the field of the one-atom maser pumped by two-level atoms injected in a coherent superposition of their upper and lower states (Slosser, Meystre and Wright, 1990).

Unfortunately, the measurement of the nonclassical photon statistics in the cavity is not that straightforward. The measurement process of the field invokes the coupling to a measuring device, with losses leading inevitably to a destruction of the nonclassical properties. The ultimate technique to obtain information about the field employs the Rydberg atoms themselves: Measure the photon statistics via the dynamic behavior of the atoms in the radiation field, i.e. via the collapse and the revivals of the Rabi oscillations, that is one possibility. However, since the photon statistics depends on the interaction time which has to be changed when collapse and revivals are measured, it is much better to probe the population of the atoms in the upper and lower maser levels when they leave the cavity. In this case, the interaction time is kept constant. Moreover, this measurement is relatively easy since electric fields can be used to perform selective ionization of the atoms. The detection sensitivity is sufficient so that the atomic statistics can be investigated. This technique maps the photon statistics of the field inside the cavity via the atomic statistics.

In this way, the number of maser photons can be inferred from the number of atoms detected in the lower level (Meschede, Walther and Müller, 1985). In addition, the variance of the photon number distribution can be deduced from the number fluctuations of the lower-level atoms (Rempe and Walther, 1990). In the experiment, we are therefore mainly interested in the atoms in the lower maser level. Experiments carried out along these lines are described in the following section.

4. Experimental Results - A Beam of Atoms with Sub-Poissonian Statistics

Under steady state conditions, the photon statistics of the field is essentially determined by the dimensionless parameter $\Theta = (N_{ex}+1)^{1/2}\Omega t_{int}$, which can be understood as a pump parameter for the one-atom maser (Filipowicz, Javanainen and Meystre, 1986). Here, N_{ex} is the average number of atoms that enter the cavity during the lifetime of the field T_c, t_{int} the time of flight of the atoms through the cavity, and Ω the atom-field coupling constant (one-photon Rabi frequency). The one-atom maser threshold is reached for Θ=1. At this value and also at $\Theta=2\pi$ and integer multiples thereof, the photon statistics is super-Poissonian. At these points, the maser field undergoes first-order phase transitions (Filipowicz, Javanainen and Meystre, 1986). In the regions between these points sub-Poissonian statistics is expected. The experimental investigation of the photon number fluctuation is the subject of the following discussion.

In the experiments (Rempe, Schmidt-Kaler and Walther, 1990), the number N of atoms in the lower maser level is counted for a fixed time interval T roughly equal to the storage time T_c of the photons. By repeating this measurement many times the probability distribution p(N) of finding N atoms in the lower level is obtained. The normalized variance (Mandel, 1979) $Q_a = [\langle N^2\rangle - \langle N\rangle^2 - \langle N\rangle]/\langle N\rangle$ is evaluated and is used to characterize the deviation from Poissonian statistics. A negative (positive) Q_a value indicates sub-Poissonian (super-Poissonian) statistics, while Q_a=0 corresponds to a Poisson distribution with $\langle N^2\rangle - \langle N\rangle^2 = \langle N\rangle$. The atomic Q_a is related to the normalized variance Q_f of the photon number by the formula

$$Q_a \quad = \quad \epsilon P Q_f(2+Q_f), \tag{1}$$

which was derived by Rempe and Walther (1990) with P denoting the probability of finding an atom in the lower maser level. It follows from formula (1) that the nonclassical photon statistics can be observed via sub-Poissonian atomic statistics. The detection efficiency ϵ for the Rydberg atoms reduces the sub-Poissonian character of the experimental result. The detection efficiency was 10% in our experiment; this includes the natural decay of the Rydberg states between the cavity and field ionization. It was determined by both monitoring the power-broadened resonance line as a function of flux (Meschede, Walther and Müller, 1985) and observing the Rabi-oscillation for constant flux but different atom-field interaction times (Rempe, Walther and Klein, 1987). In addition, this result is consistent with all other measurements described in the following, especially with those on the second maser phase transition.

Experimental results for the transition $63p_{3/2} \leftrightarrow 61d_{3/2}$ are shown in Figure 4. The measured normalized variance Q_a is plotted as a function of the flux of atoms. The atom-field interaction time is fixed at $t_{int}=50\mu s$. The atom-field coupling constant Ω is rather small for this transition, Ω=10kHz. A relatively high flux of atoms N_{ex} >10 is therefore needed to drive the one-atom maser above threshold. The large positive Q_a observed in the experiment proves the large intensity fluctuations at the onset of maser oscillation at Θ=1. The solid line is plotted according to Eq.(1) using the theoretical predictions for Q_f of the photon statistics (Filipowicz, Javanainen and Meystre, 1986; Lugiato, Scully and Walther, 1987). The error in the signal follows from the statistics of the counting distribution p(N). About 2×10^4 measurement intervals are needed to keep the error of Q_a below 1%. The statistics of the atomic beam is measured with a detuned cavity. The result is a Poisson distribution. The error bars of the flux follow from this measurement. The agreement between theory and experiment is good.

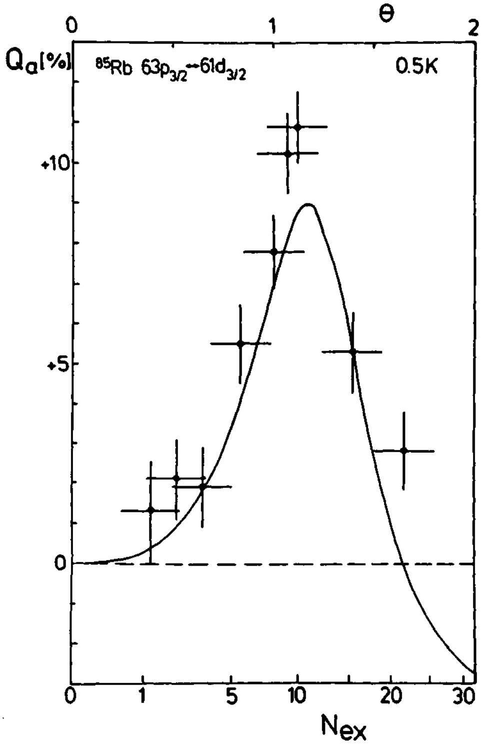
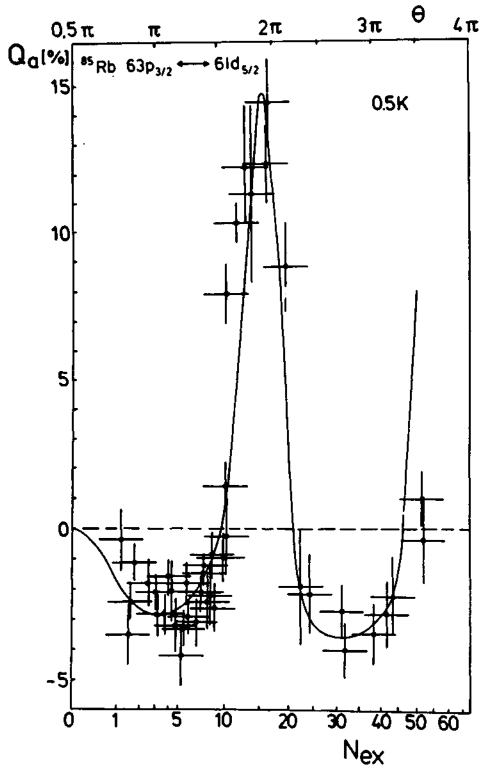


Fig. 4. Variance Q_a of the atoms in the lower maser level as a function of flux N_{ex} near the onset of maser oscillation for the $63p_{3/2} \leftrightarrow 61d_{3/2}$ transition (Rempe, Schmidt-Kaler and Walther 1990).

The nonclassical photon statistics of the one-atom maser is observed at a higher flux of atoms or a larger atom-field coupling constant. The $63p_{3/2} \leftrightarrow 61d_{5/2}$ maser transition with Ω=44kHz is therefore studied. Experimental results are shown in Figure 5. Fast atoms with an atom-cavity interaction time of t_{int}=35μs are used. A very low flux of atoms of $N_{ex} > 1$ is already sufficient to generate a nonclassical maser field. This is the case since the vacuum field initiates a transition of the atom to the lower maser level, thus driving the maser above threshold.

The sub-Poissonian statistics can be understood from Figure 6, where the probability of finding the atom in the upper level is plotted as a function of the atomic flux. The oscillation observed is closely related to the Rabi nutation induced by the maser field. The solid curve was calculated according to the one-atom maser theory with a velocity dispersion of 4%. A higher flux generally leads to a higher photon number, but for $N_{ex} < 10$ the probability of finding the atom in the lower level decreases. An increase in the photon number is therefore counterbalanced by the fact that the probability of photon emission in the cavity is reduced. This negative feed-back leads to a stabilization of the photon number (Rempe and Walther, 1990). The feed-back changes sign at a flux $N_{ex} \approx 10$, where the second maser phase transition is observed at $\Theta = 2\pi$. This is again characterized by

Fig. 5. Same as Figure 3, but above threshold for the $63p_{3/2} \leftrightarrow 61d_{5/2}$ transition (Rempe, Schmidt-Kaler and Walther 1990).



large fluctuations of the photon number. Here the probability of finding an atom in the lower level increases with increasing flux. For even higher fluxes, the state of the field is again highly nonclassical. The solid line in Figure 5 represents the result of the one-atom maser theory using Eq.(1) to calculate Q_a. The agreement with the experiment is very good. The sub-Poissonian statistics of atoms near N_{ex}=30, Q_a=-4% and P=0.45 (see Figure 6) is generated by a photon field with a variance $\langle n^2 \rangle - \langle n \rangle^2 = 0.3 * \langle n \rangle$, which is 70% below the shot noise level. Again, this result agrees with the prediction of the theory (Filipowicz, Javanainen and Meystre, 1986; Lugiato, Scully and Walther, 1987). The mean number of photons in the cavity is about 2 and 13 in the regions $N_{ex} \approx 3$ and $N_{ex} \approx 30$, respectively. Near $N_{ex} \approx 15$, the photon number changes abruptly between these two values. The next maser phase transition with a super-Poissonian photon number distribution occurs above $N_{ex} \approx 50$.

Sub-Poissonian statistics is closely related to the phenomenon of antibunching, for which the probability of detecting a next event shows a minimum immediately after a triggering event. The duration of the time interval with reduced probability is of the order of the coherence time of the radiation field. In our case this time is determined by the storage time of the photons. The Q_a value therefore depends on the measuring interval T. The

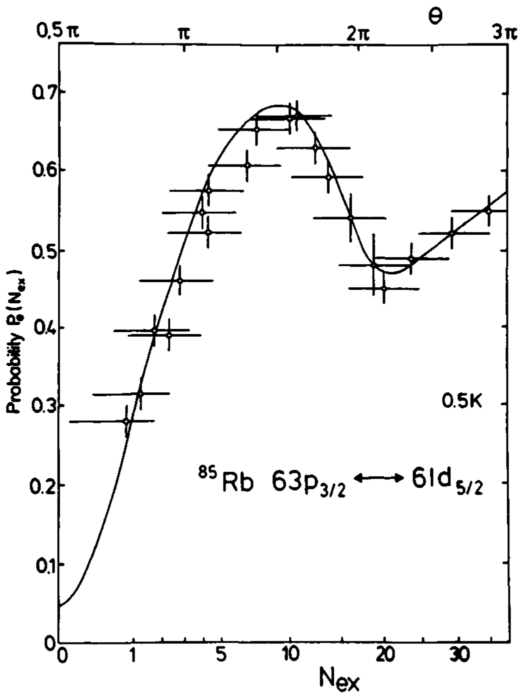


Fig. 6. Probability $P_e(N_{ex})$ of finding the atom in the upper maser level $63p_{3/2}$ for the $63p_{3/2} \leftrightarrow 61d_{5/2}$ transition as a function of the atomic flux.

measured Q_a value approaches a time-independent value for $T>T_c$. For very short sampling intervals, the statistics of atoms in the lower level shows a Poisson distribution. This means that the cavity cannot stabilize the flux of atoms in the lower level on a time scale which is short in relation to the intrinsic cavity damping time.

We want to emphasize that the reason for the sub-Poissonian atomic statistics is the following: A changing flux of atoms changes the Rabi-frequency via the stored photon number in the cavity. By adjusting the interaction time, the phase of the Rabi-nutation cycle can be chosen so that the probability for the atoms leaving the cavity in the upper maser level increases when the flux and therefore the photon number is enlarged or vice versa. We observe sub-Poissonian atomic statistics in the case where the number of atoms in the lower state is decreasing with increasing flux and photon number in the cavity. The same argument can be applied to understand the nonclassical photon statistics of the maser field: Any deviation of the number of light quanta from its mean value is counterbalanced by a correspondingly changed probability of photon emission for the atoms. This effect leads to a natural stabilization of the maser intensity by a feed-back loop incorporated into the dynamics of the coupled atom-field system.

The experimental results presented here clearly show the sub-Poissonian photon statistics of the one-atom maser field. An increase in the flux of atoms leads to the predicted second maser phase transition. In addition, the maser experiment leads to an atomic beam with atoms in the lower maser level showing number fluctuations which are up to 40% below those of a Poissonian distribution found in usual atomic beams. This is interesting, because atoms in the lower level have emitted a photon to compensate for cavity losses inevitably present under steady-state conditions. But this is a purely dissipative phenomenon giving rise to fluctuations. Nevertheless the atoms still obey sub-Poissonian statistics.

5. A New Probe of Complementarity in Quantum Mechanics

The preceding section discussed how to generate a nonclassical field inside the maser cavity. But this field is extremely fragile because any attenuation causes a considerable broadening of the photon number distribution. Therefore it is difficult to couple the field out of the cavity while preserving its nonclassical character. But what is the use of such a field? In the present section we want to propose a new series of experiments performed inside the maser cavity to test the "wave-particle" duality of nature, or better said "complementarity" in quantum mechanics.

Complementarity (Bohm, 1951; Jammer, 1974) lies at the heart of quantum mechanics: Matter sometimes displays wave-like properties manifesting themselves in interference phenomena, and at other times it displays particle-like behavior thus providing "which-path" information. No other experiment illustrates this wave-particle duality in a more striking way than the classic Young's double-slit experiment (Wootters and Zurek, 1979; Wheeler and Zurek, 1983). Here we find it impossible to tell which slit light went through while observing an interference pattern. In other words, any attempt to gain which-path information disturbs the light so as to wash out the interference fringes. This point has been emphasized by Bohr in his rebuttal to Einstein's ingenious proposal of using recoiling slits (Wheeler and Zurek, 1983) to obtain "which-path" information while still observing interference. The physical positions of the recoiling slits, Bohr argues, are only known to within the uncertainty principle. This error contributes a random phase shift to the light beams which destroys the interference pattern.

Such random-phase arguments illustrating in a vivid way how the "which-path" information destroys the coherent-wave-like interference aspects of a given experimental setup, are appealing. Unfortunately, they are incomplete: In principle, and in practice, it is possible to design experiments which provide "which-path" information via detectors which do not disturb the system in any noticeable way. Such "Welcher Weg"- (German for "which-path") detectors have been recently considered within the context of studies involving spin coherence (Englert, Schwinger and Scully, 1988). In the present section we describe a quantum optical experiment (Scully and Walther, 1989) which shows that the loss of coherence occasioned by "Welcher Weg" information, that is, by the presence of a "Welcher Weg"-detector, is due to the establishing of quantum correlations. It is in no way associated with large random-phase factors as in Einstein's recoiling slits.

The details of this application of the micromaser are discussed by Scully, Englert and Walther (1991). Here only the essential features are given. We consider an atomic interferometer where the two particle beams pass through two maser cavities before they reach the two slits of the Youngs interferometer. The interference pattern observed is then also determined by the state of the maser cavity. The interference term is given by:

$$\langle \Phi_1^{(f)}, \Phi_2^{(i)} \mid \Phi_1^{(i)}, \Phi_2^{(-f)} \rangle,$$

where $\mid \Phi_j^{(i)} \rangle$ and $\mid \Phi_j^{(f)} \rangle$ denote the initial and final states of the maser cavity.

Let us prepare, for example, both one-atom masers in coherent states $| \Phi_j^{(i)} \rangle = | \alpha_j \rangle$ of large average photon number $\langle n \rangle = | \alpha_j |^2 \gg 1$. The Poissonian photon number distribution of such a coherent state is very broad, $\Delta n \approx \alpha \gg 1$. Hence the two fields are not changed much by the addition of a single photon associated with the two corresponding transitions. We may therefore write

$$| \Phi_j^{(f)} \rangle \cong | \alpha_j \rangle$$

which to a very good approximation yields

$$\langle \Phi_1^{(f)}, \Phi_2^{(i)} | \Phi_1^{(i)}, \Phi_2^{(-f)} \rangle, \cong \langle \alpha_1, \alpha_2 | \alpha_1, \alpha_2 \rangle = 1.$$

Thus there is an interference cross term different from zero. When we, however, prepare both maser fields in number states $|n_j\rangle$ (Krause, Scully and Walther, 1987; Krause, Scully, Walther and Walther, 1989; Meystre, 1987; Meystre, Rempe and Walther, 1988) the situation is quite different. After the transition of an atom to the d-state, that is after emitting a photon in the cavity the final states read

$$| \Phi_j^{(f)} \rangle = | n_j + 1 \rangle$$

and hence

$$\langle \Phi_1^{(f)}, \Phi_2^{(i)} | \Phi_1^{(i)}, \Phi_2^{(f)} \rangle = \langle n_1, n_2 | n_1, n_2 + 1 \rangle = 0,$$

that is the coherence cross term vanishes and no interferences are observed.

On first sight this result might seem a bit surprising when we recall that in the case of a coherent state the transition did not destroy the coherent cross term, i.e. did not affect the interference fringes. However, in the example of number states we can, by simply "looking" at the one-atom maser state, tell which "path" the atom took.

It should be pointed out that the beats disappear not only for a number state. For example, a thermal field leads to the same result. In this regard, we note that it is not enough to have an indeterminate photon number to ensure interferences. The state $| \Phi_j^{(f)} \rangle$ goes as $a_j^\dagger \, | \Phi_j^{(i)} \rangle$ where $a_j^\dagger$ is the creation operator for the j-th maser. Hence the inner product

$$\langle \Phi_j^{(i)} | \Phi_j^{(f)} \rangle \to \langle \Phi_j^{(i)} | a_j^\dagger | \Phi_j^{(i)} \rangle,$$

and in terms of a more general density matrix formalism we have

$$\langle \Phi^{(i)} | \Phi^{(f)} \rangle \to \sum_n \sqrt{n+1} \rho_{n,n+1}^{(i)}.$$

Thus we see that an off-diagonal density matrix is needed for the production of beats. For example, a thermal field having indeterminate photon number would not lead to interferences since the photon number distribution is diagonal in this case.

The atomic interference experiment in connection with one-atom maser cavities is a rather complicated scheme for a "Welcher Weg"-detector. There is a much simpler possibility which we will discuss briefly in the following. This is based on the logic of the famous "Ramsey fringe" experiment. In this experiment two microwave fields are applied to the atoms one after the other. The interference occurs since the transition from an upper state to a lower state may either occur in the first or in the second interaction region. In order to

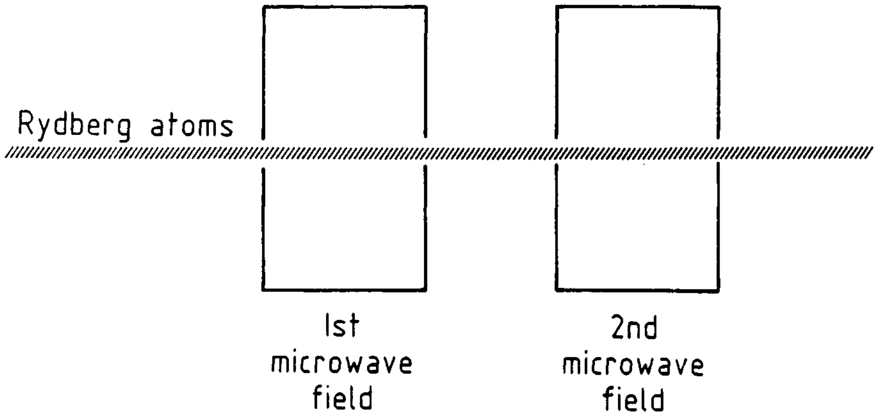


Fig. 7. Setup for the Ramsey experiment.

calculate the transition probability we must sum the two amplitudes and then square, thus leading to an interference term. We will show here only the principle of this experiment; a more detailed discussion is the subject of another paper (Englert, Walther and Scully, 1992). In the setup we discuss here, the two Ramsey fields are two one-atom maser cavities (see Figure 7). The atoms enter the first cavity in the upper state and are weakly driven into a lower state $|b\rangle$. That is, each microwave cavity induces a small transition amplitude governed by $m\tau$ where m is the atom-field coupling constant and τ is the time of flight across the cavity.

Now if the quantum state of the initial (final) field in the j-th cavity is given by $\Phi_j^{(i)}$ ($\Phi_j^{(f)}$) then the state of the atom + maser 1 + maser 2 at the various relevant times is given in terms of the coupling constant m_j and interaction times τ_j, initial $|\Phi_j^{(i)}\rangle$ and final states $|\Phi_j^{(f)}\rangle$ of the j-th maser by:

$$
\begin{aligned}
|\psi(0)\rangle\!\rangle &= |a, \Phi_1^{(i)}, \Phi_2^{(i)}\rangle \\
|\psi(\tau_1)\rangle\!\rangle &\cong |a, \Phi_1^{(i)}, \Phi_2^{(i)}\rangle - im_1\tau_1 \;\; |b, \Phi_1^{(f)}, \Phi_2^{(i)}\rangle \\
|\psi(\tau_1 + T)\rangle\!\rangle &\cong |a, \Phi_1^{(i)}, \Phi_2^{(i)}\rangle - im_1\tau_1 \;\; |b, \Phi_1^{(f)}, \Phi_2^{(i)}\rangle e^{-i\Delta\omega T} \\
|\psi(\tau_1 + T + \tau_2)\rangle\!\rangle &\cong |a, \Phi_1^{(i)}, \Phi_2^{(i)}\rangle - im_1\tau_1 \;\; |b, \Phi_1^{(f)}, \Phi_2^{(i)}\rangle e^{-i\Delta\omega T} \\
&\quad -im_2\tau_2 \;\; |b, \Phi_1^{(i)}, \Phi_2^{(f)}\rangle
\end{aligned}
$$

where $\Delta\omega$ is the atom-cavity detuning and T >> τ_j the time of flight between the two cavities. If we ask for P_b, the probability that the atom exits cavity 2 in the lower state $|b\rangle$, this is given by

$$P = \left[\langle\Phi_1^{(f)},\Phi_2^{(i)}\,|m_1^*\tau_1 e^{i\Delta\omega \mathrm{T}} + \langle\Phi_1^{(i)},\Phi_2^{(f)}\,|m_2^*\tau_2\right]$$
$$\times\left[|\Phi_1^{(f)},\Phi_2^{(i)}\rangle m_1\tau_1 e^{-i\Delta\omega \mathrm{T}} + |\Phi_1^{(i)},\Phi_2^{(f)}\rangle m_2\tau_2\right]$$
$$= m_1^* m_1\tau_1^2 + m_2^* m_2\tau_2^2 + (m_1^* m_2\tau_1\tau_2 e^{i\Delta\omega \mathrm{T}}\langle\Phi_1^{(f)},\Phi_2^{(i)}|\Phi_1^{(i)},\Phi_2^{(f)}\rangle + c.c.).$$

Now in the usual Ramsey experiment $|\,\Phi_j^{(i)}\rangle = |\,\Phi_j^{(f)}\rangle = |\,\alpha_j\rangle$, where $|\,\alpha_j\rangle$ is the coherent state in the j-th maser, which is not changed by the addition of a single photon. Thus the "fringes" appear going as $\exp(i\Delta\omega\mathrm{T})$. However, consider the situation in which $|\,\Phi_j^{(i)}\rangle$ is a number state e.g. the state $|\,0_j^{(i)}\rangle$ having no photons in the j-th cavity initially; now we have

$$P = m_1^* m_1\tau_1^2 + m_2^* m_2\tau_2^2 + (m_1^* m_2\tau_1\tau_2 e^{i\Delta\omega \mathrm{T}}\langle 1_1, 0_2\,|\,0_1, 1_2\rangle \;+\; c.c.).$$

In this case, the one-atom masers are now acting as "Welcher Weg"-detectors, and the interference term vanishes due to the atom-maser quantum correlation.

We note that the more usual Ramsey fringe experiment involves a strong field "$\frac{\pi}{2}$-pulse" interaction in the two regions. This treatment is more involved than necessary for the present purposes. A more detailed analysis of the one-atom maser Ramsey problem is given elsewhere (Englert, Walther and Scully, 1992).

We conclude this section by emphasizing again that this new and potentially experimental example of wave-particle duality and observation in quantum mechanics displays a feature which makes it distinctly different from the Bohr-Einstein recoiling-slit experiment. In the latter the coherence, that is the interference, is lost due to a phase disturbance of the light beams. In the present case, however, the loss of coherence is due to the correlation established between the system and the one-atom maser. Random-phase arguments never enter the discussion. We emphasize that the argument of the number state not having a well-defined phase is not relevant here; the important dynamics is due to the atomic transition. It is the fact that which-path information is made available which washes out the interference cross terms (Scully, Englert and Walther, 1991).

6. Conclusion

In this paper a review of the work with the one-atom-maser was given emphasizing the tests of the Jaynes-Cummings model. The experiments show nicely that this model, originally considered as a treatment of a purely academic problem, got nice confirmation after so many years.

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