Axion detection in the Micromaser

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Abstract

We report on a scheme for highly efficient detection of single microwave photons, with application to detecting certain exotic particles such as the axion. This scheme utilises an experiment known as the micromaser and the phenomenon of trapping states to amplify the signal and produce a cascade of detector counts in a field ionization detector. The cascade provides a time-resolved signal marking the presence of a single photon. For experimentally achievable parameters, the detection efficiency exceeds 90%.

Key words: Micromaser; Rydberg atoms; Cavity QED; Axion; Photon detection; Trapping states; Quantum trajectories

Reliable rare-event detection is vital in many fields of physics, especially in particle physics where there is a real need to observe the existence of rare or weakly interacting particles. One such particle that has received much attention is the axion \( \text{m} \), which is predicted to have a mass equivalent to a microwave photon. While it is very difficult to detect the axion directly schemes do exist for converting these particles into microwave photons via the Primakoff process \( \text{2,3} \), where the problem becomes one of detecting the subsequent microwave photon. The use of Rydberg atoms in single microwave photon detectors \( \text{4,5} \) is an attractive prospect for this purpose since it is experimentally easier to detect the effect of the photons on atoms than the photons themselves. The typical method for detecting whether a Rydberg atom has absorbed or emitted a microwave photon is using “state selective field ionization” \( \text{6} \), where the atoms are subjected to a varying electric field that ionizes them in a position dependent upon their state. Detecting a single atom in an altered state reveals the presence of a photon. However, the nature of these detectors is such that a number of effects will produce detector clicks that have nothing to do with a signal and are unavoidable in practice. This has important practical implications when trying to resolve a single detection event \( \text{5} \) in the presence of background thermal noise.
The scheme proposed in this paper produces a signal that can be discriminated from single clicks arising from random clicks unconnected with a signal, as the arrival of a single photon triggers a cascade of detector clicks. The microwave photon detection efficiency approaches 100% and is also inherently more robust to dark counts (counts in the absence of signal), missed counts (due to finite detector efficiency) and mis-counts (arising from detector cross talk). This method of detection is based on a quantum mechanical effect called zero photon population trapping in the micromaser (7, 8, 9). The micromaser plays the role of an ultrasensitive microwave detector for axions created in an ancillary conversion chamber (5, 10). Mode matching between the two cavities would provide efficient coupling to any type of axion conversion chamber and therefore we will not comment on the structure of the conversion chamber itself. We will instead concentrate this discussion on the structure of the detector itself and describe how it achieves a high detection efficiency.

The micromaser (11) is a cavity QED experiment in which we use superconducting microwave cavities with a $Q$-factors up to $5 \times 10^{10}$ (single photon lifetime of around 0.4s), through which we pass a sequence of two level atoms that interact one at a time with a single mode of the cavity. The atoms are prepared initially in their excited state and the transition between the two states is resonant with the cavity. The atoms are very strongly coupled to the cavity mode, therefore allowing a maser field to be produced with only one atom passing through the cavity at a time. While the microwave photons themselves cannot be detected, the atoms can, so we are able to derive a great deal of information about the field from the atoms emerging from the cavity. The micromaser has already been used to observe the appearance of single quanta and as a triggerable source of single microwave quanta (12) and can be used as a microwave photon detector via population trapping states. They occur when the cavity state is trapped and emission of the incident atoms is forbidden until the arrival (by some other means) of a single photon, causing a cascade of emission events. Thus the single photon is massively amplified by the micromaser. This occurs regardless of the atomic pump rate, which can be up to several thousand atoms per second. When a single photon enters the cavity the conditions change from destructive to constructive interference and emission probability can rise to nearly 100%. Thus the rate of emission events can go from zero to thousands per second on the arrival of one photon.

Figure 1 shows the experimental operation of the micromaser. Rubidium atoms are emitted from an oven in their ground state (unfilled circles) in a highly collimated beam. A laser excites these atoms to the $63P_{3/2}$ Rydberg state (black circles), which acts as the upper level $|e\rangle$ of what is effectively a two level system. In this case, the lower level $|g\rangle$ is the $61D_{5/2}$ Rydberg state, separated by 21.5GHz. However, Rydberg transitions are closely spaced and span frequencies from 10–120GHz, allowing us to search anywhere in a mass range of approximately 40–500µeV. The excitation laser is angled with respect
Fig. 1. Schematic of operation of the micromaser. Ground state rubidium atoms (small unfilled circle) exit the oven with thermal velocities. A detuned angled laser excitation region excites a particular velocity class to the $63P_{3/2}$ Rydberg state (black circles). The transition between this and the $61D_{5/2}$ state is resonant with a single mode in a superconducting microwave cavity and interacts coherently with it. The atomic state is recorded by state selective field ionization detectors upon exiting the cavity.

to the atomic beam to allow velocity selection via Doppler detuning. A typical velocity resolution of 0.5% is achievable with current techniques. The excited atoms enter the high-$Q$ superconducting cavity and interact resonantly with a single mode (typically the TE$_{121}$ mode) of the resonator. This interaction is (to a very good approximation) described by the Jaynes-Cummings Hamiltonian (13)

$$\hat{H} = \hbar\omega_0\hat{\sigma}_z + \hbar\omega \left( \hat{a}^\dagger \hat{a} + \frac{1}{2} \right) + \hbar g \left( \hat{a}^\dagger \hat{\sigma}_+ + \hat{a} \hat{\sigma}_- \right)$$

where $\omega_0$ and $\omega$ are the atomic transition and field mode frequencies respectively, $\hat{\sigma}_z$ is the atomic projection operator, $\hat{\sigma}_\pm$ are the atomic raising and lowering operators, $g$ is the coupling strength ($\approx 40\text{krads}^{-1}$ in the micromaser) between atom and field, $\hat{a}$ and $\hat{a}^\dagger$ are the photon annihilation and creation operators. This is one of the simplest Hamiltonians in quantum optics, describing the interaction of a single two-level atom with a single field mode. There are no loss mechanisms in this ideal model, but since the interaction time is 3–4 orders of magnitude smaller than the time scales for losses (the micromaser cavity has a $Q$-factor of up to $5 \times 10^{10}$), this is an excellent approximation.

During the interaction, the system undergoes Rabi oscillations between the states $|e\rangle$ and $|g\rangle$. After exiting the cavity, the state of the atom is measured by state selective field ionization, giving us information about the field. The probability that an atom emits a photon into the cavity is given by

$$P_{\text{emit}} = \sin^2 \left( g\tau \sqrt{n + 1} \right)$$

where $\tau$ is the interaction time and $n$ is the number of photons already in the cavity. By tuning the atomic velocity correctly, we can reduce this probability to zero, which can be understood as the system undergoing an integer number
of Rabi oscillations,

$$\tau = \frac{k\pi}{g\sqrt{n + 1}}$$  \hspace{1cm} (3)

where \(k\) is an integer. If equation 3 is fulfilled, then the cavity field is trapped with \(n\) photons and has zero probability of progressing to \(n + 1\), which is possible even for \(n = 0\) (known as the vacuum trapping state). Such trapping states have been experimentally observed \(7, 8\). By starting in the vacuum state (which is ensured by cooling the cavity to around 40mK) and appropriately tuning the velocity (and hence the interaction time) according to the value for the vacuum trapping state, emission is forbidden and atoms are unable to emit photons. However, if \(n = 1\), the probability of emission at the same velocity is close to 93%. Thus the addition of one photon has a dramatic effect on emission probability and the count rate in the \(|g\rangle\) detector goes from zero to a detectably high number, indicating the arrival of a single photon in the cavity. Subsequent emission events change the emission probability via equation 2. These probabilities mean that the arrival of a single photon in the cavity by any mechanism causes the system to jump past the vacuum trapping state, giving a cascade of \(|g\rangle\) detector counts. This is the signature of a single photon arriving in the cavity.

The proposed experimental configuration is similar to that in \(5\), with the detection cavity replaced by the micromaser system presented here.

To examine the effectiveness of this scheme, we have employed a Quantum Trajectory Analysis method to simulate detection records of the system. This technique involves stochastically evolving a wavefunction using a combination of a non-Hermitian Hamiltonian and a set of “jump operators”. Quantum trajectory analysis is often used to calculate an approximation to the density matrix of a system, but here individual trajectories are used to simulate detection records from the experiment. The implementation of the method here is based upon that in \(14\). The non-Hermitian effective Hamiltonian is given by

$$\hat{H}_{\text{eff}} = -\frac{1}{2}i\hbar\gamma \left[(n_t + 1)\hat{a}^\dagger\hat{a} + n_t\hat{a}\hat{a}^\dagger\right] - \frac{1}{2}i\hbar R + \hbar\omega\hat{a}^\dagger\hat{a}$$  \hspace{1cm} (4)

where \(\gamma\) is the cavity decay constant, \(n_t\) is the thermal photon number and \(R\) is the rate at which atoms pass through the cavity.

The particular set of jump operators used are listed below:

$$\hat{C}_{-1} = \sqrt{\gamma (n_t + 1)}\hat{a}$$  \hspace{1cm} (5)
is the operator that represents a photon being lost to the reservoir,

$$\hat{C}_0 = \sqrt{R} \cos \left( g\tau \sqrt{n + 1} \right)$$

(6)

represents an atom traversing the cavity and exiting in its original excited state,

$$\hat{C}_1 = \sqrt{R} \frac{\sin (g\tau \sqrt{n})}{\sqrt{n}} \hat{a}^\dagger$$

(7)

is the operator representing an atom introducing a photon into the field that contains $n$ photons, and

$$\hat{C}_2 = \sqrt{\gamma n} \hat{a}^\dagger$$

(8)

is the operator representing a photon being gained from the reservoir. However, since every operator maps pure states onto pure states, and for our purposes we always begin with the (pure) vacuum state, then we can reduce the dynamics simply to jumps occurring stochastically and the wavefunction remaining unchanged in between.

A quantum trajectory simulation of the micromaser operating in the vacuum trapping state was performed, in which the ground state detector count (the rate of occurrence of jump $\hat{C}_2$) was monitored while single photons were added to the cavity at random times. Figure 2 shows an example trajectory for the ideal case, with no deviation from perfect operating conditions, in order to illustrate the principle of the operation of the detector. It shows how the field evolves inside the cavity, along with the detector clicks we see when probing the atoms.

Here we see that, for sufficiently high atomic pump rates, once the vacuum trapping state is passed then the field very quickly reaches three photons, which also gives a zero emission probability and becomes trapped again (this is the $n = 3, k = 2$ trapping state, eq. 3). The field then proceeds to rapidly oscillate between this and the two and one photon states (due to decay from the cavity), giving rise to the high count rate. We see that adding just one photon at around $t = 1.25s$ produces a detector count rate of up to around 45 counts per second, which is easily detectable, even with imperfect detectors. Notice that, for a typical microwave frequency of around 21.5GHz, as used in current micromaser experiments, this amounts to near perfect detection of an energy change of less than 90$\mu eV$.

Figure 3 shows how the detection efficiency depends upon the pump rate $R$. This simulation was performed by introducing a single photon into the cavity
Fig. 2. A pair of graphs showing how the change in the cavity photon number effects the ground state count rate for ideal detectors. Notice that a very small change in the cavity photon number can produce a very large change in count rate.

Fig. 3. A graph showing how the detection efficiency increases for increasing values of atomic flux $R$. The error bars indicate the statistical spread of simulated results.

at a random time, and if a ground state count rate above a threshold of 10Hz was achieved within a set interval, then a successful detection was said to have occurred. This process was repeated 1000 times for each value of $R$ to give an average detector efficiency. The velocity spread and temperature were both set to zero in this case.

Figure 4 displays the threshold operation (in this case with an artificially high
Fig. 4. Plots to show how the proposed threshold system would work. Once the count rate exceeds the predefined threshold, a detection event is recorded and the cavity field is allowed to relax back to the vacuum state by switching off the excitation laser or applying a $\pi-$pulse to the excited atoms to pump the cavity with ground state atoms.

threshold of 30 counts / second for illustrative purposes). To prepare for the next count period, the field is then allowed to relax back to the vacuum state, either via free decay of the field, or more quickly by pumping with ground state atoms. This period of dead time is shown in the lower plot of fig. 4 delineated by vertical dashed lines. Applying a $\pi-$pulse to the incoming atoms in state $|e\rangle$ evolves their state to $|g\rangle$, which allows for faster pump-down to the vacuum state. Additionally, by relaxing the velocity selection, for example by using a perpendicular excitation scheme, the pump-down rate to the vacuum state can be further enhanced. This method would allow detector dead-times significantly shorter than simple cavity decay alone would permit.

To be more realistic, however, it is possible to include a number of departures from the ideal conditions in the quantum trajectory method to investigate the limits imposed on the system by these factors. For example, a more complete model of this system includes practical limitations of the system. *Dark Counts* are caused by detector clicks occurring by means other than ionization of the rubidium atoms (for example simple thermal excitation in the detector or cosmic rays passing through the detector), leading to non-zero count rates when no atoms are present, giving a Poisson distributed background level. High quality electron multipliers reduce this rate to around 3 counts per second or less. *Missed Counts* arise when an atom is not ionized at all in the field ionization region, or when the liberated electron does not reach the electron
multiplier. Detector Crosstalk occurs when an atom is ionized at the wrong detector, leading to errors in the statistics of the detected atoms.

All of these errors are incorporated into the model by means of setting the detector efficiencies $\eta_g < 1$ and $\eta_e < 1$ for the ground and excited state detectors respectively and adding a random background generated with a poissonian distribution centred at $r_b$ counts/second to simulate the dark counts and crosstalk.

Other errors in the system arise from the departure from ideal operating conditions of the micromaser. The idealised model assumes that there is no spread in interaction time $\tau$, the coupling parameter $g$ is constant and that there is never more than one atom in the cavity at any time. In practice, we find that, due to the linewidths of the velocity selecting laser and atomic transition, the interaction time has a non-zero spread. Mechanical vibrations in the system may also cause variations in the parameter $g$. Hence we replace $g$ and $\tau$ with $\phi = g\tau$, drawn from a normal distribution centred at $\phi_0 = \pi$ with spread $\Delta\phi$ to represent these effects. Perhaps the major source of error, however, is the occurrence of multi-atom events. When there is more than one atom in the cavity, equation 1 no longer holds, and the more complicated interaction has a high probability of breaking the trapping state barrier and causing an erroneous detection event. The probability of an atom contributing to a single atom event is given by $P = e^{-2R\tau}$, which gives a maximum rate of $62\text{s}^{-1}$ for a 99% probability of one atom events at our vacuum trapping state. This effect is easily included in the simulation by monitoring the time between incident atoms. If two or more are present in the cavity, then the trapping state is broken by the addition of one or two photons.
Fig. 6. Plot showing the composition of background counts. The lower portion represents the noise rate due to multi-atom events and the upper portion that of the finite velocity spread.

We now see that the detection efficiency increases with increasing pump rate, but that the background counts also increase, due to the higher probability of two-atom events and outliers in the velocity distribution that disrupt the trapping state. This clearly affects the signal to noise ratio, and by plugging in real experimental numbers for our errors we can easily predict the optimum operating conditions to maximise our signal to noise ratio.

Figure 6 shows that, for experimentally realistic parameters, multi-atom effects are the largest single source of noise. The effect of these multi-atom events can be reduced by increasing the coupling $g$, and hence reducing $\tau$, or altering the distribution of incoming atoms. In the limit of uniform atomic spacing, for example, we can in principle achieve an upper limit of $R = \tau^{-1} \approx 10^4$ (a promising method of altering the distribution is currently being investigated).

In this paper we have shown that it is possible to massively amplify the characteristic signal of a single microwave photon to a level where it is easily measurable with current detector technology. Furthermore, the theoretical model presented here can be used to decide the particular operating parameters for optimal performance.

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