Measuring the temperature of the spin via weak measurement

Yusuf Turek

School of Physics and Electronic Engineering, Xinjiang Normal University, Urumqi, Xinjiang 830054, China

E-mail: yusufu1984@hotmail.com

Received 24 October 2019, revised 23 February 2020
Accepted for publication 9 March 2020
Published 22 April 2020

Abstract
In this study, we give a new proposal to measure the temperature of the spin in a sample in a magnetic resonance force microscopy system by using postselected weak measurement, and investigating the Fisher information to estimate the precision of the scheme. We show that in a high temperature regime the temperature of the spin can be measured via a weak measurement technique with proper postselection and our scheme is able to increase the precision of temperature estimation.

Keywords: weak measurement, weak value, spin

(Some figures may appear in colour only in the online journal)

1. Introduction
Measurement is a basic concept in physics and any information of the system can be obtained from the measurement. Unlike the classical measurement, in most of the quantum measurement processes the information of the system only can be obtained or estimated by using indirect measurements. Using quantum measurement terminology, in a quantum measurement process, there have an interaction between the measuring device (pointer) and measured system. To guarantee the measurement precision, the interaction time must be very short so that the system and the probe itself does not affect the results of measurements. In mathematics, these requirements of quantum measurement can be expressed by the von Neumann Hamiltonian $H = g\hat{\Lambda} \otimes \hat{P}$ [1]. Here, $\hat{\Lambda}$ is the operator of the measured system we want to measure, $\hat{P}$ the canonical momentum of measuring device, and $g$ represents the measuring strength. If the interaction strength between the system and the pointer is strong, i.e. $g \gg 1$, we can get the information of the system by single trial with very small error, and the Stern–Gerlah experiment is a typical example. In contrast, if the interaction strength between the pointer and the system is very small, i.e. $g \ll 1$, there have still interference between different eigenvalues of the system observable we want to measure. The latter one is called weak measurement, and after a single trial we cannot obtain the information of the system precisely.

The weak measurement, as a generalized von Neumann quantum measurement theory, was proposed by Aharonov, Albert, and Vaidman in 1988 [2], and opens a new route to measure the system information in weak coupling measurement problems. In the weak measurement technique, the coupling between the pointer and the measured systems is sufficiently weak and the obtained information by single trial is trivial. However, in this measurement theory the measurement process does not destroy the initial state of the measured system and allows us to take measurements many times, and after many trials one can statistically get enough information of the system as precisely as possible [3]. One of the distinguished properties of weak measurement compared with strong measurement is that its induced weak value of the observable on the measured system can be beyond the usual range of the eigenvalues of that observable [4]. The feature of weak value is usually referred to as an amplification effect for weak signals rather than a conventional quantum measurement and is used to amplify much weak but useful information in physical systems [5–7]. So far, the weak measurement technique has been applied in different fields to investigate very tiny effects, such as beam deflection [8–13], frequency shifts [14], phase shifts [15], angular shifts [16], velocity shifts [17], and even temperature shift [18]. Furthermore, it has been applied to solve some fundamental problems in quantum physics such as quantum paradoxes (Hardy’s paradox [19–21] and the three-box paradox [22]), quantum...
correlation and quantum dynamics [23–28], quantum state tomography [29–35], violation of the generalized Leggett–Garg inequalities [36–41] and the violation of the initial Heisenberg measurement-disturbance relationship [42, 43], etc.

We know that the temperature is a basic concept in thermodynamics and most of the properties of a matter is directly related to its temperature. According to the thermodynamics, the temperature is an intensive quantity and independent to sample’s volume and mass. Thus, it is possible to measure the temperature of nanoscale objects which are inserted in a large sample. The motion of nanoscale objects obey the quantum mechanics, and with the recent progress to manipulation of individual quantum systems, it has been possible to get temperature readings with nanometric spatial resolution [44–46]. Quantum thermometry is also applied to precisely estimate the temperature of fermionic [47, 48] and bosonic [49, 50] hot reservoirs, micromechanical resonators [51–53] and nuclear spins [54]. The nanoscale system is too fragile and measuring the temperature of such systems requires that the measurement process should not disturb them too much while keeping the high measurement precision. Furthermore, as investigated in previous studies [6], the weak measurement almost does not destroy the measured system and can get the desired system information via statistical techniques with high precision. Thus, here raising an intriguing question as to whether one can measure the temperature of nanoscale objects by using a postselected weak measurement technique. In a recent study [55], they investigated the method to measure the temperature of a bath using the postselected weak measurement scheme with a finite dimensional probe, and anticipated more applications of postselected weak measurement method in thermometry.

In this paper, motivated by the previous study Pati et al [55], we investigate measuring the temperature of a nanoscale system (spin) which is inserted in a large thermal reservoir by using postselected weak measurement. We assume that the spin and cantilever are considered as a measured system and measuring device (pointer), respectively, and have a weak interaction between the spin and cantilever in a magnetic resonance force microscopy (MRFM) system [56]. We show that the temperature of the spin, which stayed in hot reservoir with temperature $T$, can be estimated by the postselected weak measurement technique with proper postselected states of the spin, and the weak value can be read out from optical experimental processes. In the temperature estimation process, the fluctuation always exists due to the indirect measurement, but the quantum estimation theory provides the tool to evaluate lower bounds to the amount of fluctuations for a given measurement. To investigate the precision of our method, we evaluate the Fisher information (FI) for the estimation of temperature via postselected weak measurement. We found that in a high temperature regime the FI is much larger than unity with proper postselected states of the spin, and this enable us to show that the postselected weak measurement method indeed can be used in the temperature estimation process in a MRFM system with high precision.

The rest of the paper is organized as follows. In section 2, we give the setup for our system. In section 3, we give the details how to measure the temperature of a spin via postselected weak measurement. In section 4, we investigate the Fisher information to estimate the precision of our scheme. We give a conclusion to our paper in section 5.

2. Model setup

As illustrated in figure 1, two ends of a cantilever used in MRFM is fixed and a small ferromagnetic particle is attached at the middle of the cantilever [56]. The force produced by a single spin on the ferromagnetic particle affects the parameters of the mechanical vibrations of the cantilever. There is an external permanent magnetic field $B_{\text{ext}}$ directed in the $z$ direction exerted on the spin and a ferromagnetic particle attached on the middle of the cantilever also exerts a gradient dipole magnetic field produced by the ferromagnetic particle on the spin $\partial B_d/\partial z$ with $B_d = \frac{2}{3} \mu_0 M_0 \left( \frac{R_0}{R_0 + d + z^2} \right)$. Here $\mu_0$ is the permeability of the free space, $d$ is the distance between the bottom of the ferromagnetic particle and the spin is initially at equilibrium position $(z = 0)$, $R_0$ and $M_0$ are the radius and magnetization of the ferromagnetic particle, respectively. The spin is also exposed to the transversal frequency modulated rf magnetic field $B_0(t) = B_0 \cos(\omega t + \phi(t))$, $\sin(\omega t + \phi(t))$, 0) in the $x - y$ plane. The motion of the cantilever with the ferromagnetic particle can be modeled as a simple harmonic oscillator with effective mass $m$ and frequency $\omega_1$. We can express the oscillation position $z$ of the cantilever with creation and annihilation operators $\hat{a}$ and $\hat{a}^\dagger$ by $\hat{z} = \sigma (\hat{a} + \hat{a}^\dagger)$, with $\sigma = \sqrt{\hbar / 2m\omega_1}$. In the rotating system coordinates (RSC), the Hamiltonian of this system can be written as [56]

$$H = \hbar \omega_1 \hat{a}^\dagger \hat{a} + \hbar \omega_2 \hat{S}_z + \hbar \omega_{2g} \hat{S}_z - \hbar g \hat{S}_z \otimes \hat{z}. \quad (1)$$

Here, $\omega_R = \gamma B_1$, $\omega_2 = \gamma B_0 - \omega - \frac{\partial \phi}{\partial \omega}$, $g = \gamma \frac{\partial B_d}{\partial z}$, $\gamma$ is the gyromagnetic ratio and $B_0 = B_{\text{ext}} + B_0^{(0)}$ is the total magnetic field on the spin when $z = 0$ in $z$ direction, and the gradient of the dipole field $\partial B_d/\partial z$ is taken at the spin location when
3. Measure the temperature of the spin

As given in equation (1), the interaction Hamiltonian between the system (spin) and pointer (cantilever) is given by the standard von Neumann Hamiltonian [4]

\[ H_{\text{ini}} = -g \hat{S}_z \otimes \hat{\mathbf{2}}. \]  

(2)

Here, \( g \) is a coupling coefficient between the spin and cantilever, and its value depends on the gyromagnetic ratio \( \gamma \) and the gradient magnetic field at the spin position produced by the ferromagnetic particle. \( \hat{\mathbf{2}} \) is the position operator of the pointer, while the conjugate momentum operator is \( \hat{p} = \int \hat{p} |p\rangle \langle p| dp \) with \( [\hat{p}, \hat{z}] = i\hbar \).

We consider the spin as a system and cantilever as a measuring device (pointer) to measure the temperature of the spin. We assume that initially the system (spin) is in a heat bath with temperature \( T \), and reached the thermal equilibrium state \( \rho_t = e^{-\beta H_t} / \text{Tr}(e^{-\beta H_t}) \).

Here, \( \beta = 1 / k_B T \), \( k_B \) is Boltzmann constant and it is taken as unity hereafter. We assume that the initial state of the measuring device and system are prepared as \( \phi(z) = \langle z | \phi \rangle \) and \( \rho(0) = \rho_t \), respectively. Under the action of the unitary evolution operator \( U(t) = \exp(-i \int_0^t H_{\text{ini}}(\tau) d\tau) \), the initial state of total system \( \rho(0) = \rho_t \otimes |\phi\rangle \langle \phi| \) will be evolved to

\[ \rho(t) = U^\dagger(t) \rho(0) U(t). \]  

(3)

This is the state of the total system after interaction, but we are only interested in the final state of the measuring device (cantilever) which contains system information. If we assume that the interaction strength \( g \) is very small, then it is enough to consider the expansion of the unitary operator up to its first order, and equation (3) becomes

\[ \rho(t) \approx \rho_t \otimes |\phi\rangle \langle \phi| + igt \hat{S}_z \otimes \hat{\mathbf{2}}, \quad \rho_t \otimes |\phi\rangle \langle \phi|. \]  

(4)

According to the standard terminology and procedure of weak measurement theory, if we choose the state \( |\psi\rangle \) as the final state of the spin, and project it onto the total system state, equation (4), it gives the unnormalized final state of measuring device as

\[ \phi(t) \approx \langle \psi | \rho_t | \psi \rangle e^{igt \hat{S}_z \otimes \hat{\mathbf{2}}} |\phi\rangle e^{-igt \hat{S}_z \otimes \hat{\mathbf{2}}}. \]  

(5)

Here,

\[ S_n = \frac{\langle \psi | \hat{S}_z | \psi \rangle}{\langle \psi | \rho_t | \psi \rangle} \]  

is the weak value of spin \( z \) component and is related to the initial state of the spin which reached thermal equilibrium with a heat bath with temperature \( T \). As can see from equation (5) and equation (6), we can deduce the spin temperature by properly choosing the final state \( |\psi\rangle \). However, we have to note that \( |\psi\rangle \) cannot be the eigenstate of operator \( \hat{S}_z \), otherwise the temperature will not be related to the weak value and our scheme will lose its validity to measure the temperature of the spin.

To readout the temperature of the spin system, next we will study the property of the weak value of spin operator \( \hat{S}_z \) by assuming that it is in the hot bath with high temperature, i.e. \( \beta \ll 1 \). Furthermore, we assume that the eigenvalues of \( \hat{S}_z \) and \( \hat{S}_x \) are \( \epsilon_m \) and \( \epsilon_m \), and the corresponding eigenstates are \( |\psi_m\rangle \) and \( |\psi_m\rangle \), i.e., \( \hat{S}_z |\psi_m\rangle = \epsilon_m |\psi_m\rangle \) and \( \hat{S}_x |\psi_m\rangle = \epsilon_m |\psi_m\rangle \), respectively. Then, the weak value of the spin \( z \) component can be rewritten as

\[ S_n = \frac{\langle \psi | \hat{S}_z e^{-i\beta H_t} | \psi \rangle}{\langle \psi | e^{-i\beta H_t} | \psi \rangle} \]

\[ = \frac{\sum_m e^{-\beta(\epsilon_m + \epsilon_m)} \langle \psi_m | \hat{S}_z | \psi_m \rangle \langle \psi_m | \psi_m \rangle}{\sum_m e^{-\beta(\epsilon_m + \epsilon_m)} \langle \psi_m | \psi_m \rangle \langle \psi_m | \psi_m \rangle} \]

\[ = \frac{\langle \psi | \hat{S}_z | \psi \rangle}{\langle \psi | \rho_t | \psi \rangle} - \beta \langle \psi | \hat{S}_x \hat{H}_t | \psi \rangle \frac{1}{1 - \beta \langle \psi | \hat{H}_t | \psi \rangle} \].

(7)

During the derivation of equation (7), we use \( e^{\pm x} \approx 1 \pm x \), \( (x \ll 1) \), and completeness of the bases \( |\psi_m\rangle \) and \( |\psi_m\rangle \), i.e., \( \sum_m |\psi_m\rangle \langle \psi_m | = \sum_m |\phi_m\rangle \langle \phi_m | = 1 \). Since \( (1 - x)^{-1} \approx 1 - x \), \( (x \ll 1) \), we can rewrite the equation (7) as below

\[ S_n = \frac{\langle \psi | \hat{S}_z | \psi \rangle}{\langle \psi | \rho_t | \psi \rangle} - \beta \langle \psi | \hat{S}_x \hat{H}_t | \psi \rangle \langle 1 - \beta \langle \psi | \hat{H}_t | \psi \rangle \rangle. \]  

(8)

From this relation, we can express the temperature parameter \( \beta \) by using the weak value \( S_n \), as

\[ \beta = \frac{S_n - \langle \hat{S}_z \rangle}{\langle \psi | \hat{S}_z \hat{H}_t | \psi \rangle - \langle \psi | \rho_t | \psi \rangle \langle 1 - \beta \langle \psi | \hat{H}_t | \psi \rangle \rangle}, \]  

(9)

where \( \text{Conv}(S_n, H_t) = \langle \psi | \hat{S}_z \hat{H}_t | \psi \rangle - \langle \psi | \rho_t | \psi \rangle \langle 1 - \beta \langle \psi | \hat{H}_t | \psi \rangle \rangle \). From equation (9) we can see that since the value of \( \langle \hat{S}_z \rangle \) and \( \text{Conv}(S_n, H_t) \) can be found by using the postselected state \( |\psi\rangle \), if we can find the value of \( S_n \), then the temperature of the spin can be estimated very easily. For example, if we take \( |\psi\rangle = \frac{1}{\sqrt{2}} (|\psi_1\rangle + |\psi_2\rangle) \), then equation (9) is reduced to

\[ \beta = \frac{2(1 - 2S_n)}{\omega_R + 3\omega_z}. \]  

(10)

From the above theoretical results we can deduce that the estimation of the temperature of spin in the MRFM system depends on the postselected weak value. Thus, in the remaining part of this section we will discuss the possibility of finding the weak value of our scheme. We assume that the motion of the cantilever can be described by a simple
harmonic oscillator, and its initial state can be written as
\[ \phi_0(z) = \langle z | 0 \rangle = \left( \frac{1}{2\pi\sigma^2} \right)^{\frac{3}{4}} \exp \left( -\frac{z^2}{4\sigma^2} \right), \]
where \( \sigma \) is the width of the oscillation beam. The normalized final state of the cantilever after postselection, i.e., equation (5), can be written as
\[ |\phi_f\rangle = \kappa |0\rangle + \hat{g}_0 |s_w\rangle |1\rangle, \]
where \( \hat{g}_0 = g \sigma, \kappa = (1 + g_s^2 |S_w|^2)^{-\frac{1}{2}} \), and \( \phi_0(z) = \langle z | 0 \rangle \) and \( \phi_1(z) = \langle z | 1 \rangle \) represent the ground and first excited states of the cantilever, respectively. The expectation values the position and momentum operators under the final state of the cantilever \( |\phi_f\rangle \) can be calculated as
\[ \langle z \rangle = -2\sigma g_0 \kappa^2 3S_w \]
and
\[ \langle p \rangle = \kappa^2 \sigma^{-1} g_0 \text{Re}(S_w), \]
respectively. Thus, the weak value of the spin \( z \) component is directly proportional to the shifted values of position and momentum operators of the pointer, respectively. According to the results of recent studies [57, 58], we can measure the real and imaginary parts of the weak values with optical experiments. Finally, in our scheme the temperature of spin can be measured in the lab with proper optical experimental setups.

4. Precision of our scheme

In this section, we will check the precision of our scheme by investigating the Fisher information. The Fisher information is the maximum amount of information about the parameter that we can estimate from the system. For a pure quantum state \( |\phi\rangle \), the quantum Fisher information estimating the parameter \( \beta \) is
\[ F(\beta) = 4 [\langle \partial_\beta |\phi\rangle \partial_\beta |\phi\rangle - \langle |\phi\rangle \partial_\beta |\phi\rangle |^2], \]
where the state \( |\phi\rangle \) is the normalized final state of the pointer, i.e., equation (12). The variance of unknown parameter \( \Delta \beta \) is bounded by the Cramér–Rao bound
\[ \text{Var}(\beta) \geq \frac{1}{N F(\beta)}, \]
where \( N \) is the number of measurements. From this definition of variance of parameter \( \beta \), we can see that the Fisher information set the minimal possible estimate for parameter \( \beta \), while higher Fisher information means a better estimation. As the variance of \( \beta \) is inverse proportional to the measurement time, we consider \( N = 1 \) throughout this paper.

We investigate the variation of quantum Fisher information for different postselected states \( |\psi_f\rangle = \cos \frac{\theta}{2} |1\rangle + e^{i\varphi} \sin \frac{\theta}{2} |1\rangle \). Here, \( \theta \in [0, \pi], \varphi \in [0, 2\pi], \) and \( |1\rangle \) and \( |\rangle \) are eigenvectors of the spin operator \( \hat{S}_z \), with corresponding eigenvalues \( \pm \frac{1}{2} \) and \( -\frac{1}{2} \), respectively. Our numerical results are shown in figure 2 and 3, respectively, and it revealed that the quantum Fisher information is higher in a high temperature regime, and the phase of the postselected state also can affect the precision of our scheme. Furthermore, as indicated in figure 2, when the postselected state is an eigenstate of the spin operator \( \hat{S}_z \), the Fisher information becomes zero and our scheme loses its validity to measure the temperature of the spin. Here, we have to mention that the experimental realization of our scheme needs a noiseless environment (prepare the cantilever into the ground state), and this requirement could be satisfied by current cooling technology, i.e., optomechanical cooling of the microresonator [59].

5. Conclusion

In this study, we have proposed a new scheme to measure the temperature of spin which put it in a thermal bath with temperature \( T \) by using a postselected weak measurement method. We found that in the high temperature regime, the precision of our proposal is high enough and can be controlled by adjusting...
the parameters of postselected states of the spin. Here, we only limited to a high temperature regime and did not consider the low temperature case. However, in [60], the authors showed that quantum Fisher information is higher in postselected measurement rather than non-postselected weak measurement. Thus, in our scheme if the postselected weak measurement is used in a noiseless environment, it will show good performance in a low temperature regime rather than non-post-selected measurement schemes.

Acknowledgments

This work was supported by the National Natural Science Fund Foundation of China (Grant No. 11865017).

ORCID iDs

Yusuf Turek @ https://orcid.org/0000-0002-0637-5271

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