Entanglement in the Quantum Heisenberg XY model

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1. INTRODUCTION

Quantum entanglement has been studied intensely in recent years due to its potential applications in quantum communication and information processing [1], such as quantum teleportation [2], superdense coding [3], quantum key distribution [4], and telecoloning [5]. Recently Dür et al. [6] found that truly tripartite pure state entanglement quantum computer [7] can be realized in quantum-Hall system [17] and in cavity QED system [18] for a quantum computer.

Here we consider the quantum Heisenberg XY model, which was intensively investigated in 1960 by Lieb, Schultz, and Mattis [15]. Recently Imamoǧlu et al have studied the quantum information processing using quantum dot spins and cavity QED [10] and obtained an effective interaction Hamiltonian between two quantum dots, which is just the XY Hamiltonian. The effective Hamiltonian can be used to construct the controlled-NOT gate [8]. The XY model is also realized in the quantum-Hall system [17] and in cavity QED system [18] for a quantum computer.

The XY Hamiltonian is given by

$$H = J \sum_{n=1}^{N} (S_x^n S_x^{n+1} + S_y^n S_y^{n+1})$$

where $S^α = σ^α/2$ ($α = x, y, z$) are spin 1/2 operators, $σ^α$ are Pauli operators, and $J > 0$ is the antiferromagnetic exchange interaction between spins. We adopt the periodic boundary condition, i.e., $S^x_{N+1} = S^x_1$, $S^y_{N+1} = S^y_1$.

One role of the XY model in quantum computation is that it can be used to construct the swap gate. The evolution operator of the corresponding two-qubit XY model is given by

$$U(t) = \exp \left[ -iJt(\sigma_x^1 \sigma_x^2 + \sigma_y^1 \sigma_y^2)/2 \right].$$

Choosing $Jt = π/2$, we have

$$U \left( \frac{\pi}{2J} \right) |00\rangle = |00\rangle, \quad U \left( \frac{\pi}{2J} \right) |11\rangle = |11\rangle,$$

$$U \left( \frac{\pi}{2J} \right) |01\rangle = -i|01\rangle, \quad U \left( \frac{\pi}{2J} \right) |10\rangle = -i|10\rangle.$$

The above equation shows that the operator $U \left( \frac{\pi}{2J} \right)$ acts as a swap gate up to a phase. Another gate $\sqrt{\text{swap}}$ which is universal can also be constructed simply as $U \left( \frac{\pi}{4J} \right)$.

The entanglement in the ground state of the Heisenberg model has been discussed by O’Connor and Wooters [20]. Here we study the entanglement in the XY model.
model. We first consider the generation of \( W \) states in the \( XY \) model. It is found that for 3 and 4 qubits, the \( W \) states can be generated at certain times. By the concept of concurrence, we study the entanglement properties in the time evolution of the \( XY \) model. Finally we discuss the thermal entanglement in the two-qubit \( XY \) model with a magnetic field and in the anisotropic \( XY \) model.

II. SOLUTION OF THE \( XY \) MODEL

With the help of raising and lowering operators \( \sigma^\pm_n = S^x_n \pm iS^y_n \), the Hamiltonian \( H \) is rewritten as (\( J = 1 \))

\[
H = \frac{1}{2} \sum_{n=1}^{N} (\sigma^+_n \sigma^-_{n+1} + \sigma^+_n \sigma^-_{n+1}) .
\]

Obviously the states with all spins down \( |0\rangle^\otimes N \) or all spins up \( |1\rangle^\otimes N \) are eigenstates with zero eigenvalues. The eigenvalue problem of the \( XY \) model can be exactly solved by the Jordan-Wigner transformation [2]. Here we are only interested in the time evolution problem and in the ‘one particle’ states (\( N - 1 \) spins down, one spin up),

\[
|k\rangle = \frac{1}{\sqrt{N}} \sum_{n=1}^{N} a_{k,n} \sigma^+_n |0\rangle^\otimes N .
\]

The eigenstate is given by

\[
H|\Psi\rangle = \frac{1}{2} \sum_{n=1}^{N} (a_{k,n+1} + a_{k,n-1}) \sigma^+_n |0\rangle^\otimes N
= E_k \sum_{n=1}^{N} a_{k,n} \sigma^+_n |0\rangle^\otimes N .
\]

Then the coefficients \( a_{k,n} \) satisfy

\[
\frac{1}{2} (a_{k,n+1} + a_{k,n-1}) = E_k a_{k,n} .
\]

The solution of the above equation is

\[
a_{k,n} = \exp \left( \frac{i2\pi nk}{N} \right) (k = 1...N) ,
\]

\[
E_k = \cos \left( \frac{2\pi k}{N} \right) ,
\]

where we have used the periodic boundary condition.

So the eigenvectors are given by

\[
|k\rangle = \frac{1}{\sqrt{N}} \sum_{n=1}^{N} \exp \left( \frac{i2\pi nk}{N} \right) \sigma^+_n |0\rangle^\otimes N
\]

which satisfy \( \langle k|k'\rangle = \delta_{kk'} \). It is interesting to see that all the eigenstates are generalized \( W \) states (Eq.\( ^2 \)).

Note that the \( XY \) Hamiltonian \( H \) commutes with the operator

\[
Q = \sigma^+_2 \otimes \sigma^-_3 \otimes ... \otimes \sigma^-_N ,
\]

then the state

\[
|k\rangle^\prime = \frac{1}{\sqrt{N}} \sum_{n=1}^{N} \exp \left( \frac{-i2\pi nk}{N} \right) |k\rangle
\]

are also the eigenstates of \( H \) with eigenvalues \( \cos (2\pi k/N) \).

Now we choose the initial state of the system as \( \sigma^+_1 |0\rangle^\otimes N \), and in terms of the eigenstates \( |k\rangle \), it can be expressed as

\[
|\Psi(0)\rangle = \frac{1}{\sqrt{N}} \sum_{k=1}^{N} \exp \left( \frac{-i2\pi k}{N} \right) |k\rangle .
\]

The state vector at time \( t \) is easily obtained as

\[
|\Psi(t)\rangle = \sum_{n=1}^{N} b_n(t) \sigma^+_n |0\rangle^\otimes N ,
\]

where

\[
b_n(t) = \frac{1}{\sqrt{N}} \sum_{k=1}^{N} e^{i2\pi(n-1)k/N - it \cos(2\pi k/N)} .
\]

If we choose the initial state as \( \sigma^-_1 |1\rangle^\otimes N \), then the wave vector at time \( t \) will be \( \sum_{n=1}^{N} b_n(t) \sigma^-_n |1\rangle^\otimes N \).

III. GENERATION OF \( W \) STATES

From Eq.\( ^{10} \), the probabilities at time \( t \) for state \( \sigma^+_n |0\rangle^\otimes N \) is obtained as

\[
P(n, N, t) = |b_n(t)|^2 .
\]

For \( N = 2 \), it is easy to see that the probability \( P(1, 2, t) = \cos^2 t \), \( P(2, 2, t) = \sin^2 t \). The state vector at time \( t \) is

\[
|\Psi(t)\rangle = \cos t |01\rangle - i \sin t |10\rangle .
\]

When \( t = \pi/4 \), the above state is the maximally entangled state.

Now we consider the case \( N = 3 \). The probabilities are analytically obtained as

\[
P(1, 3, t) = \frac{1}{9} \left[ 5 + 4 \cos \left( \frac{3}{2} t \right) \right] ,
\]

\[
P(2, 3, t) = P(3, 3, t) = \frac{1}{9} \left[ 2 - 2 \cos \left( \frac{3}{2} t \right) \right] .
\]
Fig. 1(a) gives a plot of the probabilities versus time. It is clear that there exist some cross points of the probabilities. At these special times the probabilities \( P(n, 3, t) \) are all equal to \( 1/3 \), which indicates the \( W \) states are generated. From Eq. (20), we see that if the time \( t \) satisfies the equation
\[
\cos \left( \frac{3}{4} t \right) = -\frac{1}{2},
\]
the probabilities are same. The solution of Eq. (21) is
\[
t_n = \frac{4\pi}{9} + \frac{4n\pi}{3}, \quad t'_n = \frac{8\pi}{9} + \frac{4n\pi}{3} (n = 0, 1, 2, ...).
\]
Explicitly at these time points, the corresponding state vectors are
\[
\left| \Psi(t_n) \right\rangle = \frac{1}{\sqrt{3}} \left( |100\rangle + e^{-i\frac{2\pi}{3}} |010\rangle + e^{-i\frac{2\pi}{3}} |001\rangle \right),
\]
\[
\left| \Psi(t'_n) \right\rangle = \frac{1}{\sqrt{3}} \left( |100\rangle + e^{i\frac{2\pi}{3}} |010\rangle + e^{i\frac{2\pi}{3}} |001\rangle \right).
\]
which are the generalized \( W \) state for \( N = 3 \).

As seen from Fig. 2(a), there also exists some cross points, which indicates the 4-qubit \( W \) states are generated. The probabilities are same when
\[
t_n = \frac{\pi}{2} + 2n\pi, \quad t'_n = \frac{3\pi}{2} + 2n\pi (n = 0, 1, 2, ...)
\]
Explicitly the 4-qubit \( W \) states are
\[
\left| \Psi(t_n) \right\rangle = \frac{1}{2} \left( |1000\rangle - i|0100\rangle - |0010\rangle - i|0001\rangle \right),
\]
\[
\left| \Psi(t'_n) \right\rangle = \frac{1}{2} \left( |1000\rangle + i|0100\rangle - |0010\rangle + i|0001\rangle \right).
\]

Can we generate \( W \) states for more than 4 qubits in the \( XY \) model? Fig. 3(a) shows that there is no cross points for \( N = 5 \). Further numerical calculations for long time and large \( N \) show no evidence that there exist some times at which the \( W \) states can be generated.

For the case \( N = 4 \), the probabilities are given by
\[
P(1, 4, t) = \cos^4 \left( \frac{t}{2} \right), \quad P(3, 4, t) = \sin^4 \left( \frac{t}{2} \right),
\]
\[
P(2, 4, t) = P(4, 4, t) = \frac{1}{4} \sin^2 t.
\]
defined as
\begin{equation}
C_{12} = \max \{\lambda_1 - \lambda_2 - \lambda_3 - \lambda_4, 0\}
\end{equation}
where the quantities \(\lambda_1 \geq \lambda_2 \geq \lambda_3 \geq \lambda_4\) are the square roots of the eigenvalues of the operator
\begin{equation}
\rho_{12} = \rho_{12}(\sigma_y \otimes \sigma_y)\rho_{12}^*(\sigma_y \otimes \sigma_y).
\end{equation}
The nonzero concurrence implies that the qubits 1 and 2 are entangled. The concurrence \(C_{12} = 0\) corresponds to an unentangled state and \(C_{12} = 1\) corresponds to a maximally entangled state.

We consider the entanglement in the state \(|\Psi(t)\rangle\). By direct calculations, the concurrence between any two qubits \(i\) and \(j\) are simply obtained as
\begin{equation}
C_{ij}(t) = 2|b_i(t)b_j(t)|.
\end{equation}

The numerical results for the concurrence are shown in Fig.1(b), Fig.2(b) and Fig.3(b).

For \(N = 3\), Fig.1(b) shows that the entanglement is periodic with period \(4\pi/3\). At times \(4n\pi/3 (n = 1, 2, 3, \ldots)\), the state vectors are disentangled and become the state \(|100\rangle\) up to a phase. The concurrences of \(C_{12}(t)\) and \(C_{13}(t)\) are same, and have two maximum points in one period, while the concurrence \(C_{12}(t)\) has only one maximum point. Fig.2(b) shows the concurrences for \(N = 4\). They are periodic with period \(2\pi\). In one period there are two disentanglement points, \(t = \pi, 2\pi\). For both concurrences \(C_{12}(t)\) and \(C_{23}(t)\), there are two maximum points in one period. If we choose large \(N\) (see Fig.3(b) for \(N = 5\)), there exists no exact periodicity for the entanglements of two qubits. From the time evolution of the concurrences we can see clearly when the system becomes disentangled and when the system maximally entangled.

V. THERMAL ENTANGLEMENT

Recently the concept of thermal entanglement was introduced and studied within one-dimensional isotropic Heisenberg model [23]. Here we study this kind of entanglement within both the isotropic XY model with a magnetic field and the anisotropic XY model.

A. Isotropic XY model with a magnetic field

We consider the two-qubit isotropic antiferromagnetic XY model in a constant external magnetic field \(B\),
\begin{equation}
H = B \left(\sigma_1^z + \sigma_2^z\right) + J \left(\sigma_1^+ \sigma_2^- + \sigma_2^+ \sigma_1^-\right).
\end{equation}
The eigenvalues and eigenvectors of \(H\) are easily obtained as
\begin{align}
H|00\rangle &= -B|00\rangle, \quad H|11\rangle = B|11\rangle, \\
H|\Psi^\pm\rangle &= \pm J|\Psi^\pm\rangle,
\end{align}
where $|\Psi^{\pm}\rangle = \frac{1}{\sqrt{2}}(|01\rangle \pm |10\rangle)$ are maximally entangled states.

The state of the system at thermal equilibrium is $\rho(T) = \exp\left(-\frac{H}{kT}\right)/Z$, where $Z = \text{Tr}[\exp\left(-\frac{H}{kT}\right)]$ is the partition function and $k$ is the Boltzmann’s constant. As $\rho(T)$ represents a thermal state, the entanglement in the state is called thermal entanglement [23].

In the standard basis, \{\ket{00}, \ket{01}, \ket{10}, \ket{11}\}, the density matrix $\rho(T)$ is written as ($k = 1$)

$$
\rho(T) = \frac{1}{2(\cosh \frac{J}{T} + \cosh \frac{B}{T})} \begin{pmatrix}
    e^{-\frac{J}{T}} & 0 & 0 & 0 \\
    0 & \cosh \frac{J}{T} - \sinh \frac{J}{T} & 0 & 0 \\
    0 & -\sinh \frac{J}{T} & \cosh \frac{J}{T} & 0 \\
    0 & 0 & 0 & e^{\frac{B}{T}}
\end{pmatrix}
$$

From Eqs. (27), (28) and (32), the concurrence is given by

$$
C = \max \left(\frac{\sinh \frac{J}{T} - 1}{\cosh \frac{J}{T} + \cosh \frac{B}{T}}, 0\right).
$$

Then we know $C = 0$ if $\sinh \frac{J}{T} \leq 1$, i.e., there is a critical temperature

$$
T_c = \frac{J}{\text{arcsinh}(1)} \approx 1.1346J,
$$

the entanglement vanishes for $T \geq T_c$. It is interesting to see that the critical temperature is independent on the magnetic field $B$.

For $B = 0$, the maximally entangled state $|\Psi^-\rangle$ is the ground state with eigenvalue $-J$. Then the maximum entanglement is at $T = 0$, i.e., $C = 1$. As $T$ increases, the concurrence decreases as seen from Fig.4 due to the mixing of other states with the maximally entangled state. For a high value of $B$ (say $B = 1.2$), the state $|00\rangle$ becomes the ground state, which means there is no entanglement at $T = 0$. However by increasing $T$, the maximally entangled states $|\Psi^{\pm}\rangle$ will mix with the state $|00\rangle$, which makes the entanglement increase (see Fig.4).

From Fig.5 we see that there is a evidence of phase transition for small temperature by increasing magnetic field. Now we do the limit $T \to 0$ on the concurrence, we obtain

$$
\lim_{T \to 0} C = \begin{cases} 
1 & \text{for } B < J, \\
1/2 & \text{for } B = J, \\
0 & \text{for } B > J.
\end{cases}
$$

So we can see that at $T = 0$, the entanglement vanishes as $B$ crosses the critical value $J$. This is easily understand since we see that if $B > J$, the ground state will be the unentangled state $|00\rangle$. This special point $T = 0, B = J$, at which entanglement becomes a nonanalytic function of $B$, is the point of quantum phase transition [24].
The density matrix $\rho(T)$ in the standard basis is given by

$$\rho(T) = \frac{1}{2coshJ_\gamma} \begin{pmatrix} csh\frac{J_\gamma}{T} & 0 & 0 & -sinh\frac{J_\gamma}{T} \\ 0 & csh\frac{J_\gamma}{T} - sinh\frac{J_\gamma}{T} & 0 & 0 \\ -sinh\frac{J_\gamma}{T} & 0 & cosh\frac{J_\gamma}{T} & 0 \\ 0 & 0 & 0 & cosh\frac{J_\gamma}{T} \end{pmatrix}$$

(37)

The square root of the eigenvalues of the operator $\varrho_{12}$ are $\frac{csh\frac{J_\gamma}{T} + cosh\frac{J_\gamma}{T}}{2cosh\frac{J_\gamma}{T}}$ and $\frac{csh\frac{J_\gamma}{T} - sinh\frac{J_\gamma}{T}}{2cosh\frac{J_\gamma}{T}}$. Then from Eq.(27), the concurrence is given by

$$C = \max\left(\frac{sinh\frac{J_\gamma}{T} - cosh\frac{J_\gamma}{T}}{cosh\frac{J_\gamma}{T} + cosh\frac{J_\gamma}{T}}, 0\right)$$

(38)

As we expected Eq. (38) reduces to Eq. (33) with $B = 0$ when $\gamma = 0$. When $\gamma = 1$, the concurrence $C = 0$, which indicates that no thermal entanglement appears in the two-qubit Ising model. In this anisotropic model, the concurrences are the same for both positive $J$ and negative $J$, i.e., the thermal entanglement is the same for the antiferromagnetic and ferromagnetic cases. The critical temperature $T_c$ is determined by the nonlinear equation

$$\sinh\frac{J_\gamma}{T} = \cosh\frac{J_\gamma}{T}$$

which can be solved numerically.

In Fig.6 we give a plot of the concurrence as a function of temperature $T$ for different anisotropic parameters. At zero temperature the concurrence is 1 since no matter what the sign of $J$ is and what the values of $\gamma$ are, the ground state is one of the Bell states, the maximally entangled state. The concurrence monotonically decreases with the increase of temperature until it reaches the critical value of $T$ and becomes zero. The numerical calculations also show that the critical temperature decreases as the anisotropic parameter increases from 0 to 1.
VI. CONCLUSIONS

In conclusion, we have presented some interesting results in the simple \(XY\) model. First, we can use \(XY\) interaction to generate the 3-qubit and 4-qubit \(W\) entangled states. Second, we see that the time evolution of entanglement are periodic for 2, 3, 4 and 6 qubits, and there is no exact periodicity for large \(N\). At some special points the states becomes disentangled. Finally we study the thermal entanglement within a two-qubit isotropic \(XY\) model with a magnetic field and an anisotropic \(XY\) model, and find that the thermal entanglement exists for both ferromagnetic and antiferromagnetic cases. Even in the simple model we see some evidence of the quantum phase transition.

The entanglement is not completely determined by the partition function, i.e., by the usual quantum statistical physics. It is a good challenge to study the entanglement in multi-qubit quantum spin models.

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