Quantum Coherence Oscillations in Antiferromagnetic Chains

Marius Grigorescu and Mahi R. Singh

Department of Physics and Astronomy
University of Western Ontario
London, Ontario, Canada N6A 3K7

Abstract:
Macroscopic quantum coherence oscillations in mesoscopic antiferromagnets appear when the anisotropy potential creates a barrier between the antiferromagnetic states with opposite orientations of the Neel vector. This phenomenon is studied for the physical situation of the nuclear spin system of eight Xe atoms arranged on a magnetic surface along a chain. The oscillation period is calculated as a function of the chain constant. The environmental decoherence effects at finite temperature are accounted assuming a dipole coupling between the spin chain and the fluctuating magnetic field of the surface. The numerical calculations indicate that the oscillations are damped by a rate $\sim (N - 1)/\tau$, where $N$ is the number of spins and $\tau$ is the relaxation time of a single spin.

PACS numbers: 75.45.+j, 73.40.Gk, 75.30.Pd, 76.60.-k
I. Introduction

The observation of macroscopic quantum coherence (MQC) phenomena in complex many-particle systems represents a subject of wide interest, ranging from the conceptual foundations of quantum mechanics \[1\], to the physics of the microelectronic devices. During the last years a particular attention was given to the macroscopic quantum tunneling and quantum coherence oscillations \[2\]. In these phenomena the quantum dynamics of a prepared non-stationary wave-packet is directly reflected by the non-classical behavior of a macroscopic observable. However, it is difficult to identify physical situations where such coexistence of the classical and quantum aspects could be observed.

Quantum coherence oscillations may occur in the localization probability of relatively complex systems as the individual Xe atoms trapped in the surface-tip junction of the scanning tunneling microscope \[3, 4, 5, 6\]. Particularly suitable candidates to observe MQC phenomena are also the magnetic systems \[7\]. In the anisotropic antiferromagnets the Neel vector may change the orientation by quantum tunneling \[8\] or quantum coherence oscillations \[9\]. The observation of these oscillations still represents a challenging problem, but important results have been obtained from the measurements of the ac magnetic susceptibility in antiferromagnetic (AF) ferritin \[10, 11\].

The study of the non-elementary excitations in AF systems is also essential for understanding the high-$T_c$ superconductivity \[12, 13\]. Therefore, a detailed study of the MQC phenomena in antiferromagnets appears highly interesting.

The purpose of this work is to investigate the occurrence of MQC oscillations of the Neel vector in an anisotropic AF chain \[C \equiv \{\vec{I}_i, i = 1, 8\}\] of eight spins \(1/2\). The Hamiltonian correspond to the system of nuclear spins for eight Xe atoms placed on a magnetic surface and coupled by the magnetic dipole interaction. Such chain structures could be constructed, for instance, using the scanning tunneling microscope in the ”atomic switch” operation mode, proved during the last years to be an efficient instrument to displace in a controlled way the Xe atoms on a Ni surface \[14\].

The model Hamiltonian and the spectrum of the spin chain are presented in Sect. II. It is found that the first two levels are quasi-degenerate, and dominated by AF states with opposite orientations of the Neel vector. Therefore, non-stationary AF wave packets which are interchanged periodically by MQC oscillations can be constructed. Their oscillation period is calculated as a function of the chain constant.

The relaxation effects at finite temperature due to the magnetic coupling between the spin chain and the surface are discussed in Sect. III. The coupling between the electronic and the nuclear magnetic moments is significant, and in the opposite situation, when the nuclear spins play the role of the environment, the MQC oscillations in small magnetic particles can be suppressed.
Here the environment is represented by the surface, and the coupling is reflected by the finite relaxation time $\tau$ of the nuclear spins. The effects of the spin relaxation on the nuclear MQC oscillations will be accounted by the coupling between the chain and the fluctuating magnetic field of the magnetic surface. It is found that when $\tau$ is in the range of seconds the MQC oscillations are damped, but not completely suppressed. The conclusions are summarized in Sect. IV.

II. Quantum coherence oscillations in the AF nuclear spin chain

The chain considered in the present calculations consists of $N = 8$ spins $1/2$ equally spaced by $d$, and interacting by the model Hamiltonian

$$H_0 = J \sum_{i=1}^{N-1} [c_x \vec{I}_i \cdot \vec{I}_{i+1} + c_y \vec{I}_i \cdot \vec{I}_{i+1} + c_z \vec{I}_i \cdot \vec{I}_{i+1}] .$$

Here $J > 0$ is the coupling strength and $c_x$, $c_y$ and $c_z$ are anisotropy coefficients. Classically, the AF ordering is described by using two sub-chains, $C_o$ and $C_e$, containing the odd, $\{\vec{I}_i, i = 1, 3, 5, ..., 2N - 1\}$ and the even spins $\{\vec{I}_i, i = 2, 4, 6, ..., 2N\}$, respectively. The corresponding magnetization vectors are $\vec{M}_o = \gamma h \sum_{i=odd} \vec{I}_i$ and $\vec{M}_e = \gamma h \sum_{i=even} \vec{I}_i$, where $\gamma$ denotes the gyromagnetic factor. The equilibrium configuration of this system is antiferromagnetic when $\vec{M}_o$ and $\vec{M}_e$ have the same magnitude $M_0$, but a relative antiparallel orientation. The energy in this case is a function of the Neel vector $\vec{n} = (\vec{M}_o - \vec{M}_e)/2M_0$, and is expressed by

$$E^A(\vec{n}) = -J(N - 1)(c_x n_x^2 + c_y n_y^2 + c_z n_z^2)/4 .$$

An anisotropic system with $c_{xy} < c_z$ has two degenerate minima, $E^A_{\text{min}} = -J(N - 1)c_z/4$ attained when $\vec{n}$ has the two possible orientations along the Z-axis, $\vec{n} = \pm \vec{e}_z$. These minima are separated by a two-dimensional potential barrier with the maximum $E^A_{\text{max}}(\phi) = -J(N - 1)(c_x \cos^2 \phi + c_y \sin^2 \phi)/4$, in the X-Y plane, where $\phi$ denotes the angle between the Neel vector and the X-axis.

A physical situation of interest described by the anisotropic Hamiltonian of Eq. (1) appears when $C = \{\vec{I}_i, i = 1, 8\}$ consists of the nuclear spins for a chain of Xe atoms placed on a magnetic surface. Such chains can be constructed, for instance, using techniques of atomic manipulations. The interaction between the Xe nuclear spins is due to the magnetic dipole forces, and the effect of the surface can be taken into account using the simple method of images. In a coordinate system with the Z-axis normal to the surface and
the Y-axis along the chain, each magnetic moment $\vec{m} = (m_x, m_y, m_z)$ lying above the surface at the distance $z = r_0 = 2.17$ Å, equal to the radius of the Xe atom, has an image $\vec{m}' = (-m_x, -m_y, m_z)$ at $z = -r_0$. The nearest neighbors of $\vec{m}$ interact also with $\vec{m}'$, and the Hamiltonian of the whole system is expressed by Eq. (1), where $J = \mu_0 h^2 \gamma^2/(4\pi d^3)$, $\mu_0 = 4\pi \times 10^{-7} N/A^2$ is the vacuum permeability, $\gamma = -1.54 \mu N/\hbar$ is the gyromagnetic factor of the $^{129}$Xe isotope and $d$ is the chain constant. The anisotropy coefficients are

$$c_x = 1 - 2 \sin^3 \alpha$$
$$c_y = -2 - 2 \sin^3 \alpha (1 - 3 \sin^2 \alpha)$$
$$c_z = 1 + 2 \sin^3 \alpha (1 - 3 \cos^2 \alpha)$$

with $\tan \alpha = d/2r_0$.

The spectrum of $H_0$ was calculated for $6.8 \text{ Å} < d < 9.4 \text{ Å}$ by solving numerically the eigenvalue equation

$$H_0 |\psi_n> = E_n |\psi_n> .$$

The basis was defined by the common eigenstates for the $Z$-components of all the spin operators $I_{i,z}$, $i=1,8$. These basis states are denoted by $|k> = |m_1^k, m_2^k, m_3^k, m_4^k, m_5^k, m_6^k, m_7^k, m_8^k>$, $k = 1, 2^8$, with $m_i^k = \pm 1/2$. The eigenstates of $H_0$ have the general form $|\psi_n> = \sum_k y_k^n |k>$. It is found that the ground ($n = 0$), and the first excited state ($n = 1$), have the largest overlap with only two basis states, denoted $|\uparrow>$ and $|\downarrow>$, which are antiferromagnetic in the classical sense. Explicitly these basis states are

$$|\uparrow> = \left| \frac{1}{2}, -\frac{1}{2}, -\frac{1}{2}, \frac{1}{2}, -\frac{1}{2}, -\frac{1}{2}, -\frac{1}{2}, -\frac{1}{2} \right>, \tag{5}$$

and

$$|\downarrow> = \left| -\frac{1}{2}, -\frac{1}{2}, -\frac{1}{2}, \frac{1}{2}, -\frac{1}{2}, -\frac{1}{2}, -\frac{1}{2}, -\frac{1}{2} \right>, \tag{6}$$

and are eigenstates of $Z$-component of the Neel operator $\vec{N} = \sum_{i=1,4}(\vec{I}_{2i-1} - \vec{I}_{2i})/4$ with eigenvalues +1 and -1.

The results will be presented in detail for the particular case of $d = 7$ Å, when $J = 0.17h/\text{ms}$ and the anisotropy coefficients given by Eq. (3) are $c_x = -0.22$, $c_y = -0.58$ and $c_z = 1.2$. The overlap coefficients between the eigenstates $|\psi_0>$ and $|\psi_1>$ and the AF states of Eq. (5) and (6) are $<\uparrow|\psi_0> = <\downarrow|\psi_0> = 0.57$, $-<\uparrow|\psi_1> = <\downarrow|\psi_1> = 0.63$. Thus, the two AF states exhaust more than 65% of the eigenstates norm. The eigenvalues $E_0 = -0.406h/\text{ms}$ and $E_1 = -0.395h/\text{ms}$ are separated by $\Delta = E_1 - E_0 =$
0.011\hbar/\text{ms} which is sensibly smaller than $E_2 - E_1 = 0.053\hbar/\text{ms}$. The small value of $\Delta$ shows that the system has a quasi-degenerate ground state. This appears as a ”tunneling doublet” determined by the two-dimensional potential barrier $E_{\text{max}}^A(\phi) = (0.066 \cos^2 \phi + 0.17 \sin^2 \phi)\hbar/\text{ms}$ separating the AF energy minima of $E_{\text{min}}^A = -0.353\hbar/\text{ms}$.

The tunneling behavior is shown clearly by the evolution of the non-stationary wave-packets prepared at the ground-state energy with a well-defined AF configuration. States with these properties are represented by the linear combinations

$$|\psi_\downarrow> = \frac{1}{\sqrt{2}}(|\psi_0 > + |\psi_1 >) \quad \text{and} \quad |\psi_\uparrow> = \frac{1}{\sqrt{2}}(|\psi_0 > - |\psi_1 >).$$

(7)

of the eigenstates $|\psi_0 >$ and $|\psi_1 >$. These wave-packets can be distinguished macroscopically by the expectation value of $N_z$, and during the time-evolution are interchanged periodically, by MQC oscillations. Thus, if the system is prepared at $t=0$ in the state $|\psi_\downarrow>$, then at the moment $t$ it will be found in the state $|\psi_\uparrow>$ with the probability $P_\uparrow(t) = \sin^2(\pi t/2T_{\text{max}})$, where $T_{\text{max}} = \hbar\pi/\Delta$. The half-period of oscillation $T_{\text{max}}$ is represented as a function of the chain constant $d$ in Fig. 1 (A). For $d = 7$ Å, one obtains $T_{\text{max}} = 0.3$ s, and the calculation of the expectation values

$$<N_k>(t) = <\psi_\downarrow|e^{iH_0t/\hbar}N_k e^{-iH_0t/\hbar}|\psi_\downarrow>$$

(8)

shows that in time $<N_x>(t) = 0$, $<N_y>(t) = 0$, while the Z-component has MQC oscillations $<N_z>(t) = <\psi_0|N_z|\psi_1> \cos(\pi t/T_{\text{max}})$, with $<\psi_0|N_z|\psi_1> = -0.83$ (Fig. 1 (B)).

It is important to emphasize the extreme sensitivity of the MQC resonance with respect to the preparation of the initial state. The expected value of the energy in the classical antiferromagnetic states, $<\downarrow|H_0|\downarrow>$ and $<\uparrow|H_0|\uparrow>$, is $E_{\text{min}}^A = -0.353\hbar/\text{ms}$, higher than $E_1 = -0.39\hbar/\text{ms}$ by $\approx 4\Delta$. This energy is far outside the interval $[E_0, E_1]$, and therefore between the classical antiferromagnetic states $|\downarrow>$ and $|\uparrow>$ there are no MQC oscillations.

### III. MQC oscillations at finite temperature

The evolution of the Neel vector at the MQC resonance, presented in the previous section, was obtained neglecting the residual interactions between the nuclear spins and the environment. However, when these interactions are considered the MQC oscillations could be damped \[17\] or completely suppressed \[18\], and the Neel vector acquires a fixed orientation. At finite temperatures the nuclear spin $\mathbf{I}_i$ of each Xe atom lying on the magnetic
surface is subject also to an external, time-dependent magnetic field $\vec{B}_e(t)$, created by the phonon modulation of the crystalline electric field and the lattice spin waves [10]. Due to this field, appears in the Hamiltonian a residual interaction term

$$H_r(t) = -\gamma \hbar \sum_{i=1,8} \vec{I}_i \cdot \vec{B}_e(t) .$$

At low temperatures the typical wave length of the surface phonons and magnons [13] is $\sim 400$ Å, much larger than the length of the chain, $L = 49$ Å. Therefore $\vec{B}_e$ will be considered to be the same for all spins, $\vec{B}_e(t) \equiv \vec{B}_e(t)$. With these approximations the residual interaction term determined by the environmental magnetic field is $H_r(t) = -\gamma \hbar \vec{B}_e(t) \cdot \vec{I}$, where $\vec{I} = \sum_{i=1,8} \vec{I}_i$.

The "atomic switch" experiments [14] have been performed at the environmental temperature $T = 4K$, when the thermal energy $k_B T = 0.34$ meV is very high compared both to the tunnel splitting $\Delta = 6.83 \cdot 10^{-12}$ meV and to the maximum barrier height $V_B = E_{\text{max}}(\pi/2) = 0.11 \cdot 10^{-9}$ meV. Therefore, the thermal environment can be considered as classical, and the field components $B^\mu_i(t), \mu = x, y, z$, will be treated as a white noise with zero mean. The normalization of this fluctuating field is ensured by the fluctuation-dissipation theorem (FDT) $< < B_\mu^\prime(t) B_\nu^{\prime *}(t') > > = \delta_{\mu, \nu} \delta(t-t')/(\gamma^2 \tau)$, where $< < ... > >$ denotes the average over the statistical ensemble describing the environment.

For a single spin of the chain, $H_r(t)$ induces transitions between the states $|1/2\rangle$ and $|-1/2\rangle$ with a rate [21] $\lambda = (| < -1/2 | I_x | 1/2 > |^2 + | < -1/2 | I_y | 1/2 > |^2 )/\tau = 1/2 \tau$, and the relaxation rate of the population difference $n_{1/2} - n_{-1/2}$ is $2\lambda = 1/\tau$. Therefore, the parameter $\tau$ has the meaning of spin-surface relaxation time, increasing as $1/T$ when $T \to 0$.

The relaxation of the MQC oscillations cannot be treated by using a similar two-level approximation, because the matrix elements of $H_r$ within the subspace generated by $|\psi_0\rangle$ and $|\psi_1\rangle$ are 0. The operators $I_x$ and $I_y$ contained in $H_r$ act on the initial state $|\psi_\perp\rangle$ by flipping the individual spins, and in time the quantum state acquires components over the whole spectrum. This process is described in principle by a transport equation for the density matrix [1, 20] but due to the relatively large number of states ($= 2^8$), such numerical calculations are not feasible. Moreover, at high temperatures, when the initial state of the system is a pure state and there are only few observable of interest, as in the present case, the computational effort required to find the evolution of the whole density matrix is not justified. Instead, an equivalent description [21] which can be applied efficiently is provided by a statistical ensemble of $N_t$ Brownian trajectories $|\psi^r(t)\rangle$, $r = 1, N_t$, obtained by integrating the Schrödinger equation

$$i\hbar \partial_t |\psi^r(t)\rangle = [H_0 - \gamma \hbar \vec{B}_e(t) \cdot \vec{I}] |\psi^r(t)\rangle .$$

(10)
The numerical integration was performed using the procedure presented in [3], as a classical system of Hamilton equations for the real and imaginary parts of the amplitudes \( y_k(t) = \langle k|\psi(t) >, k = 1,2^8 \). Using the notation \( u_k(t) \equiv Re(y_k(t)) \) and \( v_k(t) \equiv Im(y_k(t)) \), Eq. (10) takes the form

\[
2\hbar \dot{u}_k = \frac{\partial \mathcal{H}(t)}{\partial v_k}, \quad 2\hbar \dot{v}_k = -\frac{\partial \mathcal{H}(t)}{\partial u_k}, \tag{11}
\]

where

\[
\mathcal{H}(t) = \sum_{k,k'=1}^{N} (u_k u_{k'} + v_k v_{k'}) Re(\langle k|H_0 + H_r(t)|k' >) -
(u_k v_{k'} - v_k u_{k'}) Im(\langle k|H_0 + H_r(t)|k' >). \tag{12}
\]

The dominant AF states are annihilated by \( I_z \), and therefore the contribution of the term \( B^e(t)I_z \) from \( H_r \) is very small, and it was neglected. The remaining X and Y components of the fluctuating field at the moment \( t_n = n dt \), normalized according to the FDT, have the form \( B^e(t_n) = R_n \sqrt{1/(\gamma^2 \tau dt)} \) where \( \{R_n, n = 1,2,3,...\} \) is a sequence of random numbers with 0 mean and variance 1.

The time-evolution of the ensemble average for any observable \( \mathcal{O} \) is defined by

\[
\langle \langle \mathcal{O} \rangle \rangle (t) = \frac{1}{N_t} \sum_{r=1}^{N_t} \langle \psi^r(t)|\mathcal{O}|\psi^r(t) \rangle. \tag{13}
\]

The observables of interest here are the Neel vector and the energy, and their averages were calculated using Eq. (13) with \( \mathcal{O} \) denoting the components of the Neel operator \( N_k \), and \( H_0 \), respectively. The results obtained for \( d = 7 \) Å, \( \tau = 2.5 \) s and \( N_t = 20 \) are presented in Fig. (2). The components \( \langle \langle N_x \rangle \rangle \) and \( \langle \langle N_y \rangle \rangle \) of Fig. 2 (A) and (B) have thermal fluctuations around 0, while \( \langle \langle N_z \rangle \rangle \) shown in Fig. 2 (C) has damped oscillations. These oscillations can be well approximated by the analytical expression

\[
\langle \langle N_z \rangle \rangle (t) = e^{-\Lambda t} \langle N_z \rangle (t), \tag{14}
\]

where the damping constant obtained by fit is \( \Lambda = 3 \) s\(^{-1}\). Similar calculations with \( \tau \) in the range of seconds, as expected at low temperatures, indicate that \( \Lambda = 7.5/\tau \).

The ratio \( (\langle \langle H_0 \rangle \rangle (t) - E_0)/\Delta \) between the average excitation energy and the doublet splitting \( \Delta \) is a suitable measure of the heating effect produced by the environment, and is represented in Fig. 2 (D). The ensemble average of the energy appearing here is accurately reproduced by the formula

\[
\langle \langle H_0 \rangle \rangle (t) = \langle \langle H_0 \rangle \rangle (0) + \Delta (w_1 t - w_2 t^2), \quad \text{with } w_1 = 45.9 \text{ s}^{-1} \text{ and }
\]

\[
 w_2 = \frac{\Lambda}{2}. \tag{15}
\]
$w_2 = 19.3 \text{s}^{-2}$.

IV. Conclusions

In this work it was shown that the chain of nuclear spins for eight atoms of the $^{129}\text{Xe}$ isotope placed on a magnetic surface has a low-energy mode of collective excitation represented by antiferromagnetic MQC oscillations. The numerical calculations indicate that when the chain constant $d$ is larger than 6.8 Å the first excited state $\psi_1$ of the chain has a very low excitation energy. This state is strongly coupled to the ground state $\psi_0$ by the Z-component of the Neel operator, such that $| < \psi_0 | N_z | \psi_1 > | = 0.83$. Both states, $\psi_0$ and $\psi_1$, are dominated by two AF spin configurations having opposite orientations of the Neel vector. Therefore, resonant wave packets containing only one of these AF configurations can be constructed by the symmetric and antisymmetric linear combinations of $\psi_0$ and $\psi_1$. These wave-packets are non-stationary, and are interchanged periodically in time. Although the number of spins considered here is relatively small, the resonant wave packets are distinguished by the orientation of the Neel vector, and their oscillation represent a MQC phenomenon. When $d = 7$ Å the half-period of oscillation is $T_{\text{max}} = 0.3$ s, and increases exponentially with $d$ (Fig. 1 (A)).

The MQC oscillations are known to be very sensitive to the decoherence effects produced by the coupling to the thermal environment. If in practice these effects are strong, then the oscillations may be completely suppressed and cannot be observed. In the present case decoherence appears due to the residual dipole interaction between the nuclear spins and the fluctuating magnetic field produced by the surface. This coupling was described by one parameter, chosen to be the relaxation time $\tau$ of a single nuclear spin.

The relaxation rate of the MQC oscillations can be easily calculated in the physical situations which can be treated within a two-level approximation. However, for eight spins the number of states coupled by the residual dipole interaction term is large, and this approximation cannot be applied. Therefore, the fluctuating residual interaction term was included in the Hamiltonian, and the diffusion of the initial wave-packet $| \psi_\downarrow >$ was described by an ensemble of Brownian solutions obtained by integrating the Schrödinger equation. The results indicate that when $\tau$ is in the range of seconds, the MQC oscillations are not suppressed, but are damped by a rate $\sim (N-1)/\tau$, where $N$ is the number of spins.

ACKNOWLEDGMENTS: One of the authors (MS) is thankful to NSERC of Canada for financial support in the form of a research grant.
References

[1] W. H. Zurek, Physics Today 36, 37 (1992).

[2] J. A. Leggett, S. Chakravarty, A. T. Dorsey, M. P. A. Fisher, A. Garg and W. Zerger, Rev. Mod. Phys. 59, 1 (1987).

[3] M. Grigorescu, Rom. J. Phys. 42, 597 (1997).

[4] M. Grigorescu, P. Budau and N. Carjan, Phys. Rev. B55, 7244 (1997).

[5] P. Budau and M. Grigorescu, Phys. Rev. B57, 6313 (1998).

[6] I. S. Tilinin, M. A. Van Hove and M. Salmeron, Phys. Rev. B57, 4720 (1998).

[7] L. Thomas, F. Lionti, R. Ballou, D. Gatteschi, R. Sessoli, and B. Barbara, Nature (London) 383, 145 (1996).

[8] B. Barbara and E. M. Chudnowsky, Phys. Lett. 145, 205 (1990).

[9] A. Chiolero and D. Loss, Phys. Rev. Lett. 80, 169 (1998).

[10] S. Gider, D. D. Awschalom, T. Douglas, S. Mann and M. Chaparala, Science 268, 77 (1995).

[11] S. Gider, D. D. Awschalom, D. P. DiVincenzo and D. Loss, Science 272, 425 (1996).

[12] R. P. Singh and M. Singh, Phys. Rev. B46, 14069 (1992).

[13] S. Chakravarty, Science 278, 1412 (1997).

[14] D. M. Eigler, C. P. Lutz and W. E. Rudge, Nature 352, 600 (1991).

[15] A. Garg, Phys. Rev. Lett. 70, 1541 (1993), 74, 1458 (1995).

[16] C. Kittel, Introduction to Solid State Physics, Seventh Edition, John Wiley & Sons Inc. (1996).

[17] V. A. Benderskii, V. I. Goldanskii and D. E. Makarov, Phys. Rep. 233, 195 (1993).

[18] A. J. Bray and M. A. Moore, Phys. Rev. Lett. 49, 1545 (1982).

[19] M. G. Cottam and D. R. Tilley, Introduction to Surface and Superlattice Excitations, Cambridge University Press, (1989), p. 127.
[20] A. O. Caldeira, A. J. Leggett, Physica A 121, 587 (1983).

[21] M. Grigorescu, Physica A 256, 149 (1998),
http://xxx.lanl.gov/abs/quant-ph/9709033.

Figure Captions

Fig. 1. The half-period of the MQC oscillations $T_{max}$ as a function of the chain constant $d$ (Å) and the expected value $< N_z >$ as a function of time when $d = 7$ Å (B).

Fig. 2. Ensemble averages of the Neel vector components $<< N_x >>$ (A), $<< N_y >>$ (B), $<< N_z >>$ (C) and $<< H_0 >> - E_0)/\Delta$ (D) as functions of time when $d = 7$ Å and $\tau = 2.5$ s.
Fig. 2

(a) 

(b) 

(c)