# Influence of longitudinal acoustic phonons on domain wall magnetoresistance in magnetic nanowires

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(Received 21 August 2010; published 20 January 2011)

The influence of phonons on domain wall magnetoresistance in magnetic nanowires has gained crucial importance. In the present work, the resistance of a magnetic domain wall in a magnetic nanowire has been investigated in the presence of a longitudinal acoustic phonon using the semiclassical approach. The analysis has been based on the Boltzmann transport equation, within the relaxation time approximation. The resistance resulting from the one-dimensional Néel-type domain wall has been studied when the modulation of the exchange interaction by lattice vibrations and the electron-phonon interaction are considered. The results indicate that phonons play a significant role in resistance, and increasing the number of phonons leads to enhancement of the domain wall resistance. Increasing and decreasing the domain wall resistance is observed via enhancing the temperature and mass of the domain wall, respectively. The effect of phonons on the resistance of the domain wall is considerable in providing new spintronics devices based on the nanowires at high temperatures.

DOI: 10.1103/PhysRevB.83.035413 PACS number(s): 73.63.−b, 75.76.+j

## I. INTRODUCTION

Spintronics, which is based on exploiting the fact that electrons have spin as well as charge, is one of the most exciting and fastest growing fields in electronics.[1] This emerging technology promises to provide a new generation of electronic devices with increased performance and improved functionality.[2] To make such new spintronics devices possible, proper understanding of the spin transport phenomena and spin manipulation are essential.[2,3] Therefore, many studies have been conducted in order to control the spin relaxation rate and consequently the resistivity of the systems.[4] According to the results of these studies, one can realize spin manipulation by applying magnetic fields or optical pulses.[5] The most popular method for spin manipulation is using the Rashba spin-orbit interaction (SOI), because it causes the effective magnetic field that induces spin precession, and the strength of the Rashba SOI can be changed by gate bias.[6] The spin precession associated with the Rashba coupling led Datta and Das to propose a spin field-effect transistor in which the spin of an electron passing through the device is controlled by the SOI.[7] In real materials, full control of the spin relaxation rates and a good estimate of the electrical resistivity require an accurate description of the role of different scattering mechanisms such as impurities and phonons on the spin relaxation rates and electrical resistivity.[8,9] At nonzero temperatures, the study of the spin accommodation without considering the effects of electron-phonon interaction is not enough to completely understand the spin transport phenomena. Therefore, the phonon contribution to the spin relaxation rate has been investigated both theoretically and experimentally.[10–13] According to the theoretical calculations,[10,11] the phonon-induced spin flip is believed to be the mechanism behind the reduction of spin relaxation time at high temperatures, which is observed in experiments.[12,13]

In addition to the above-mentioned scattering sources, the spin-dependent transport in magnetic systems can be influenced by the geometrical orientation of the local magnetic structures such as the domain walls (DWs).[14] From a practical point of view, measuring the excess electrical resistivity caused by the DW and determining the effect of different scattering sources on the domain wall resistance (DWR) have been a topic of interest for several years.[15–20] For instance, a spin-dependent impurity was proposed by Levy and Zhang as a source for mixing the spin channels and magnetoresistance (MR) enhancement.[16] Additionally, the effect of the Rashba SOI on the DWR has been investigated.[17–19] The results indicate that SOI can decrease the spin relaxation rates and consequently increase the DWR. In recent years, applying an external magnetic field or propagating a photon to manipulate the spin of the electrons or control the DWR has been studied also.[19–21] Moreover, the temperature dependence of the resistivity of magnetic systems such as nanowires has been extensively studied both experimentally[22–26] and theoretically.[27,28] In particular, the resistivity of the Ni wires is measured at high temperatures and significant resistance enhancement has been found by increasing temperature.[22,23] Furthermore, the dominant and considerable role of the DWs in the resistivity of the Ni (Refs. 24 and 25) and Fe (Ref. 26) nanowires has been determined at room temperature. The results make clear that the DWs have a positive contribution in increasing the resistivity.[24–26] Despite this, most of the present studies have been concerned with the temperature dependence of the resistivity of magnetic materials in the absence of the DWs, whereas the temperature-dependent part of the DWR has remained largely unexplored. Therefore, a good estimate and a quantitative way of predicting the DWR associated with the electron-phonon interaction are needed and should open new insights into understanding the contribution of the DW to the resistance.

Here, we have used the semiclassical approach to study the resistance of a DW in the presence of electron-phonon interaction. The rest of this paper is organized as follows. In Sec. II, a model for describing the DWR in the presence of the phonon is first introduced, and then the scattering matrices

 

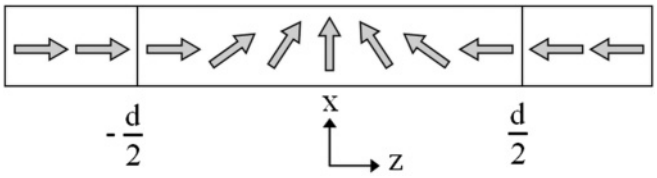


FIG. 1. A Néel-type DW with width $d$ between two magnetic regions.

and the resistance are determined. Sections III and IV are both devoted to results, discussion, and summary.

## II. THEORETICAL CONSIDERATIONS

### A. Description of model and interaction

We have studied a magnetic nanowire containing an ideal one-dimensional Néel-type DW. As shown in Fig. 1, the DW is considered between two magnetic regions with an opposite direction of magnetization. The magnetization of the DW is continuously rotated over the DW width, $d$, and it is transverse to the $z$ axis at the center of the wall. For this linear DW, the angle between the local direction of the magnetization and the $z$ axis can be written as $\theta(z) = (\pi/d)z + (\pi/2)$.

The Hamiltonian describing the DW is given by

$$H = H_{\text{el}} + H_{\text{ph}} + H_{\text{ex}} + H_{\text{el-ph}}, \tag{1}$$

where $H_{\text{el}}$ is the kinetic energy of electrons in the effective mass approximation, namely,

$$H_{\text{el}} = -\frac{\hbar^2}{2m^*}\frac{d^2}{dz^2}, \tag{2}$$

in which $m^*$ is the effective electron mass.

The second term of the Hamiltonian is the bare phonon Hamiltonian, $H_{\text{ph}}$, and can be expressed as

$$H_{\text{ph}} = \sum_q \hbar\omega_q \left(b_q^\dagger b_q + \frac{1}{2}\right); \tag{3}$$

here, $b_q^\dagger$ and $b_q$ are the phonon creation and annihilation operators that create and annihilate a phonon with wave vector $q$, respectively. It should be noted that, at sufficiently low temperatures, only acoustic phonons are relevant in metals and the longitudinal acoustic (LA) modes have the strongest coupling constants.[29] Hence, only the LA phonons have been considered in the present work. We have described a linear dispersion $\omega_q = v_s q$ that relates the frequency of a LA phonon, $\omega_q$, to its wave vector, $q$. In this equation, $v_s$ is the sound velocity and indicates the speed of the propagation of a phonon in the DW.

The third term of the Hamiltonian, $H_{\text{ex}}$, is the exchange interaction between the conduction electrons and the localized magnetic moments. In order to develop a reasonable treatment of the local onsite exchange interaction, one can consider the following scheme. The exchange interaction is assumed to be dependent on the displacement of the magnetic ions from their equilibrium sites. We have assumed that the vibrating atomic spins are spontaneously directing along the local direction of the effective mean field. It means that modulation of the local exchange interaction by the lattice vibrations is taken into account. The exchange coupling is expressed as

$$H_{\text{ex}} = -\Delta_{\text{ex}}\hat{\sigma}\cdot\hat{M}(\vec{R}+\vec{u}), \tag{4}$$

in which $\Delta_{\text{ex}}$ is the exchange interaction strength, $\hat{\sigma}$ denotes the spin operators in terms of the Pauli spin matrices, $\hat{M}(\vec{R}+\vec{u})$ is the unit vector along direction of the local magnetization, and equilibrium atomic positions are given by $\vec{R}$. The exchange interaction in terms of the displacement of magnetic ions has been expanded and only the linear terms have been kept, namely,

$$\begin{aligned} H_{\text{ex}} &= \tilde{H}_{\text{ex}}(\vec{R}) + \tilde{H}_{\text{ex}}(\vec{u}) \\ &= -\Delta_{\text{ex}}\hat{\sigma}\cdot\hat{M}(\vec{R}) - \Delta_{\text{ex}}\hat{\sigma}\cdot(\vec{u}\cdot\vec{\nabla}\hat{M}(\vec{R})), \end{aligned} \tag{5}$$

where $\vec{u}$ is the small deviation from the equilibrium position of the ion, $\vec{R}$. The Hamiltonian $\tilde{H}_{\text{ex}}(\vec{R}) = -\Delta_{\text{ex}}\hat{\sigma}\cdot\hat{M}(\vec{R})$ is now a purely electronic exchange Hamiltonian. The second term of the exchange interaction, i.e., modulated exchange, $\tilde{H}_{\text{ex}}(\vec{u})$, can be simplified to

$$\tilde{H}_{\text{ex}}(\vec{u}) = -\Delta_{\text{ex}}\left(\frac{\pi}{d}\right)[\sigma_x\cos\theta(z) - \sigma_z\sin\theta(z)](\hat{e}_q\cdot\vec{u}). \tag{6}$$

For the one-dimensional system, the displacement is determined as [30]

$$\vec{u}(z) = \sum_q \hat{e}_q\sqrt{\frac{\hbar}{2Nm\omega_q}}(b_q e^{iqz} + b_q^\dagger e^{-iqz}); \tag{7}$$

here, $N$ and $m$ are the number and the mass of magnetic ions, respectively. $\hat{e}_q$ is the unit vector that represents the displacement direction, i.e., $\hat{e}_q$ is parallel to the $z$ axis for the LA phonon propagating along the wire axis.

The last term in Eq. (1), $H_{\text{el-ph}}$, is the electron-phonon interaction Hamiltonian and can be defined by[30]

$$H_{\text{el-ph}} = D_{\text{ac}}\vec{\nabla}\cdot\vec{u}(z), \tag{8}$$

where $D_{\text{ac}}$ is the deformation potential for electron scattering by acoustic phonons. Using Eq. (7), the electron-longitudinal acoustic-phonon interaction Hamiltonian can be written as[30]

$$H_{\text{el-ph}} = D_{\text{ac}}\sum_q\sqrt{\frac{\hbar}{2Nm\omega_q}}(i\hat{e}_q\cdot\vec{q})(b_q e^{iqz} - b_q^\dagger e^{-iqz}). \tag{9}$$

By applying an approach based on the perturbation method introduced by Levy and Zhang,[16] the eigenstates of $H_{\text{el}} + \tilde{H}_{\text{ex}}(\vec{R})$ for a one-dimensional system can be determined as

$$|\psi^\uparrow(k)\rangle = \frac{\tilde{\alpha}(k)}{\sqrt{d}}e^{ikz}R_{\theta(z)}\begin{pmatrix}1\\ik\zeta\end{pmatrix}, \tag{10a}$$

$$|\psi^\downarrow(k)\rangle = \frac{\tilde{\alpha}(k)}{\sqrt{d}}e^{ikz}R_{\theta(z)}\begin{pmatrix}ik\zeta\\1\end{pmatrix}; \tag{10b}$$

here, $\tilde{\alpha}(k) = (1+k^2\zeta^2)^{-1/2}$ and $\zeta = \pi^2\hbar^2/(8m^*\Delta d^2)$ are the normalization coefficient and the perturbation parameter, respectively, in which $k$ is the electron wave vector. The perturbation parameter, $\zeta$, measures the deviation from the adiabatic transport regime.[16] Finally, $R_{\theta(z)} = \exp[-i\theta(z)\hat{\sigma}\cdot\hat{n}/2]$

is the rotation operator for electron spin, and $\hat{n}$ denotes the direction of the DW rotation axis, which in this case is assumed to be along the $y$ axis (Fig. 1). On the other hand, the eigenstate of the phonon Hamiltonian, $H_{\rm ph}$, is defined by $|n_q\rangle$, where $n_q = (e^{\hbar\omega_q/k_BT}-1)^{-1}$ denotes the Bose-Einestein distribution function of phonons. Therefore, the DW eigenstates, representing a noninteracting electron-phonon system, are $|\psi^\sigma(k),n_q\rangle = |\psi^\sigma(k)\rangle \otimes |n_q\rangle$, where $\sigma$ is the index of spinors inside the DW ($\sigma = \uparrow, \downarrow$).

### B. Spin-dependent transport characteristics

The Boltzmann equation has been employed to find the deviation from the equilibrium distribution function and to understand the physics of nonequilibrium systems. The Boltzmann equation in the relaxation time approximation reduces to

$$e\vec{E}\cdot\frac{1}{\hbar}\vec{\nabla}_k f_\sigma + \frac{1}{\hbar}\vec{\nabla}_k\varepsilon\cdot\vec{\nabla}_r f_\sigma = -\frac{f_\sigma - f_0}{\tau^\sigma}, \tag{11}$$

where $f_0$ and $f_\sigma$ are equilibrium and nonequilibrium distribution functions, respectively, $\vec{E}$ is the electric field, and $\tau^\sigma$ is the spin-dependent relaxation time. In the relaxation time approximation, this equation for a homogeneous system has a solution of the form

$$f_\sigma(k,q) = f_0(\varepsilon) - \tau^\sigma(k,q)e\vec{E}\cdot\vec{v}_{k,\sigma}\frac{\partial f_0}{\partial\varepsilon}. \tag{12}$$

Here $\vec{v}_{k,\sigma}$ is the electron velocity.[17] The spin-dependent relaxation time, $\tau^\sigma(k,q)$, can be determined by

$$[\tau^\sigma(k,q)]^{-1} = \frac{d}{\hbar}\sum_{\sigma',q'}\int dk' \left|V_s^{\sigma',\sigma}(k',k,q',q)\right|^2 \times\left[1-\frac{\tau^{\sigma'}(k',q)v_{k',\sigma'}}{\tau^\sigma(k,q)v_{k,\sigma}}\right]\delta(\varepsilon_{k\sigma}-\varepsilon_{k'\sigma'}), \tag{13}$$

in which the matrix elements of the scattering potentials are

$$V_s^{\sigma',\sigma}(k',k,q',q) = \langle\psi^{\sigma'}(k'),n_{q'}|\tilde{H}_{\rm ex}(\vec{u}) + H_{\text{el-ph}}|\psi^\sigma(k),n_q\rangle. \tag{14}$$

According to Eq. (14), we can immediately introduce $V_s^{\sigma',\sigma}(k',k,q',q)$ terms by

$$V_s^{\sigma,\sigma}(k',k,q',q) = \tilde{\alpha}(k')\tilde{\alpha}(k)\sqrt{\frac{\hbar}{2Nm\omega_q}}\Big[-\Delta_{\rm ex}\left(\frac{\pi}{d}\right) \times(k-k')\zeta + sD_{\rm ac}q(1+kk')\zeta^2\Big] \times\frac{\sin[(k-k'+sq)d/2]}{(k-k'+sq)d/2} \times\sqrt{n_q+\frac{1}{2}-\frac{s}{2}}\delta_{q',q-s}, \tag{15a}$$

$$V_s^{-\sigma,\sigma}(k',k,q',q) = \tilde{\alpha}(k')\tilde{\alpha}(k)\sqrt{\frac{\hbar}{2Nm\omega_q}} \times\left[\Delta_{\rm ex}\left(\frac{\pi}{d}\right)+sD_{\rm ac}q(k-k')\zeta\right] \times\frac{\sin\left[(k-k'+sq)d/2\right]}{(k-k'+sq)d/2} \times\sqrt{n_q+\frac{1}{2}-\frac{s}{2}}\delta_{q',q-s}, \tag{15b}$$

where $s=+1,-1$ denotes the absorption and emission of phonon inside the DW. In our study, the wire is considered in one dimension along the $z$ axis, so the integration in Eq. (13) is carried out over the one-dimensional $k$ space. Writing Eq. (13) for up and down components gives a coupled equation for $\tau^\uparrow(k,q)$ and $\tau^\downarrow(k,q)$ as

$$\tau^\uparrow(k,q) = \frac{\hbar^3}{m^*d}\left(\frac{a_{22}-a_{12}}{a_{11}a_{22}-a_{21}a_{12}}\right), \tag{16a}$$

$$\tau^\downarrow(k,q) = \frac{\hbar^3}{m^*d}\left(\frac{a_{11}-a_{21}}{a_{11}a_{22}-a_{21}a_{12}}\right), \tag{16b}$$

in which the $a_{ij}$ elements are defined as

$$a_{11} = \frac{2}{k}|V^{\uparrow\uparrow}(-k,k,q-s,q)|^2 + \frac{1}{k_-}(|V^{\downarrow\uparrow}(k_-,k,q-s,q)|^2 + |V^{\downarrow\uparrow}(-k_-,k,q-s,q)|^2), \tag{17a}$$

$$a_{12} = -\frac{1}{k}(|V^{\downarrow\uparrow}(k_-,k,q-s,q)|^2 - |V^{\downarrow\uparrow}(-k_-,k,q-s,q)|^2), \tag{17b}$$

$$a_{21} = -\frac{1}{k}(|V^{\uparrow\downarrow}(k_+,k,q-s,q)|^2 - |V^{\uparrow\downarrow}(-k_+,k,q-s,q)|^2), \tag{17c}$$

$$a_{22} = \frac{2}{k}|V^{\downarrow\downarrow}(-k,k,q-s,q)|^2 + \frac{1}{k_+}(|V^{\uparrow\downarrow}(k_+,k,q-s,q)|^2 + |V^{\uparrow\downarrow}(-k_+,k,q-s,q)|^2), \tag{17d}$$

where $k_\pm = \sqrt{k^2 \mp 4m^*\Delta/\hbar^2}$. The Fermi wave vector, $k_F$, satisfies the condition $k_F \gg \sqrt{4m^*\Delta/\hbar^2}$, and therefore the simplification $\tau^\sigma(k_\pm,q) \approx \tau^\sigma(k,q)$ for electrons with energies close to the equilibrium Fermi energy has been taken into account for deriving Eq. (16). It can be shown that the non-spin-flip forward scattering is forbidden for one-dimensional systems in the elastic regime.

Using $\tau^\sigma(k,q)$, the DW resistivity per unit length can be determined as

$$\Re = \left[\frac{e^2\hbar^2}{2\pi m^{*2}}\sum_{\sigma,q}\int dk\,k^2\tau^\sigma(k,q)\delta(\varepsilon_{k\sigma}-\varepsilon_F)\right]^{-1}, \tag{18}$$

where $\varepsilon_F$ is the Fermi energy. For evaluating the resistivity from the above equation, we have used the Debye model. In this model, the summation over $q$ is calculated by replacing it with the integral and considering the Debye wave vector, $q_D$. This vector is related to the free-electron Fermi wave vector, $k_F$, in one-dimensional metals by $q_D = \frac{2}{z}k_F$, where $z$ is the nominal valence.[31]

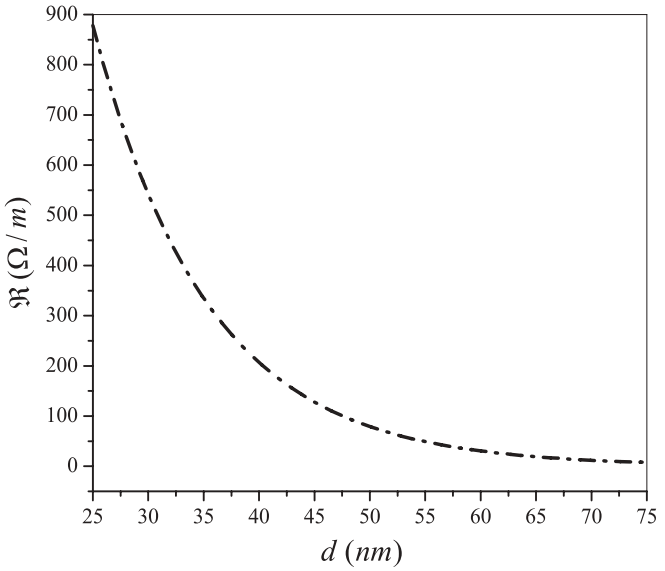


FIG. 2. DWR per unit length vs DW width at $T = 100$ K, $D_{\rm ac} = 1.0$ eV, $m = 9.746 \times 10^{-26}$ kg, and $v_s = 4970$ m/s (calculated for Ni).

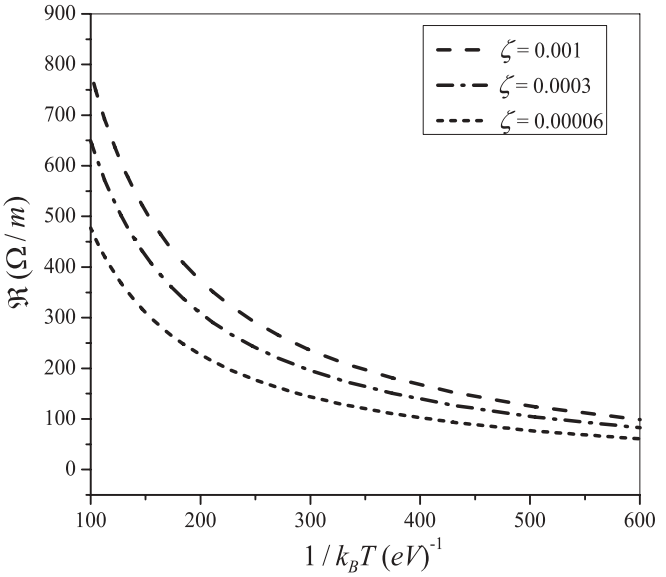


FIG. 3. DWR per unit length vs $1/k_BT$ at different nonadiabatic regimes for $D_{\rm ac} = 1.0$ eV, $m = 9.746 \times 10^{-26}$ kg, and $v_s = 4970$ m/s (calculated for Ni).

## III. RESULTS AND DISCUSSIONS

Our study has been performed on the magnetic DWs in the nanowires such as Ni and Fe. The typical parameters have been considered to be $\Delta_{\rm ex} = 0.1$–$1.0$ eV, $k_F = 17.6\ {\rm nm}^{-1}$ $(17.1\ {\rm nm}^{-1})$, $v_s = 4970$ m/s (4910 m/s) $m = 9.746 \times 10^{-26}$ kg $(9.273 \times 10^{-26}$ kg), and $m^* = m_e$, where $m_e$ is the electron mass.[31–33]

The width of the DW is the main parameter that determines the adiabaticity of the transport. Depending on this parameter, the DWs lead to a coupling of spin-up and spin-down electrons in the conduction channels. Hence, the spin adiabatically follows the local magnetization direction for very wide DWs. In addition, low-energy electrons are also transmitted adiabatically. In contrast, for narrow DWs, the spin of the carriers is not directed along the direction of the local magnetization, and the transport process is nonadiabatic.[34,35] The nonadiabaticity causes the spin channels within the DW to mix, and consequently this mixing increases the DWR. The DWR vs DW width in the presence of LA phonon is shown in Fig. 2. Our results indicate that decreasing the DW width leads to resistance enhancement. As explained, the magnitude of the resistance depends on the nonadiabaticity and mixing effect. By decreasing the DW width, these effects become more effective, so the DWR is increased.

The effect of temperature on the DWR has also been investigated. The resistance per unit length of the DW vs $1/k_BT$ at different nonadiabatic regimes is presented in Fig. 3. As the temperature increases, the number of excited phonons and consequently the effects of the exchange and the electron-phonon interactions increase. This means that increasing the number of phonons at high temperatures results in the enhancement of the DWR. Similar enhancement trends in the resistivity of Ni and Fe nanowires have been reported in the experimental studies.[22–26]

The deviation from the adiabatic regime and the nonadiabaticity in the up-spin and down-spin electron transfer through the DW depend on the $\zeta$ parameter. The nonadiabaticity causes the spin channels within the DW to mix, and consequently this mixing increases the DWR, as shown in Fig. 3.

The resistance per unit length of the DW vs $D_{\rm ac}/\Delta_{\rm ex}$ for DW with different masses of magnetic ions and sound velocities is shown in Fig. 4. The strength of the exchange and the electron-phonon interactions become more effective by increasing $\Delta_{\rm ex}$ and $D_{\rm ac}$, respectively. By enhancing $D_{\rm ac}$, the effect of the electron-phonon interaction is increased, i.e., scattering of the electron spin by the LA phonon becomes more significant. Therefore, increasing $D_{\rm ac}$ leads to the enhancement of the resistance. In contrast, increasing $\Delta_{\rm ex}$ decreases the resistance of the DW, which can be expected owing to the

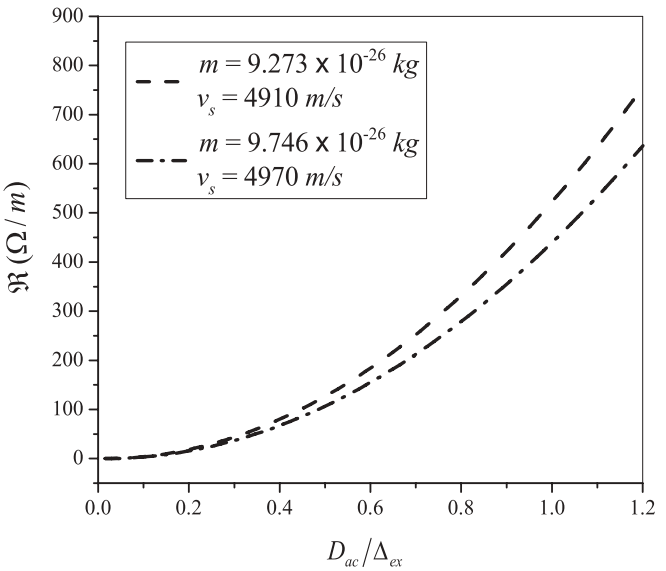


FIG. 4. DWR per unit length vs $D_{\rm ac}/\Delta_{\rm ex}$ at different masses of ions and sound velocities in the DW for $d = 40$ nm and $T = 100$ K (calculated for Ni and Fe).

fact that increasing $\Delta_{ex}$ increases the adiabatic transport ratio. Although the size of the DW is introduced as the main factor for adiabatic transport, nonadiabaticity inside the DW can be attributed to the weak magnitude of the exchange interaction, which reduces the effect of the DW on the adiabatic transport. In the case of high exchange coupling, the spin follows the local magnetization direction adiabaticity. Adiabatic transport inside the DW removes the mixing of the spin channels and decreases the resistance associated with the DW. It should be noted that the problem of adiabaticity will become more complete if other scattering mechanisms are taken into account. In the present study, increasing the strength of the electron-phonon interaction has a critical effect and enhances significantly the DWR.

In addition, Fig. 4 indicates the dependence of resistance on the mass of magnetic ions and the sound velocities of Ni and Fe. It is clear that the ionic mass has a crucial role in the resistance, and increasing the mass leads to a decrease of the DWR. This can be explained by the fact that a phonon in a DW with high ionic mass cannot produce considerable deviation from equilibrium positions of the ions and significant vibrations in the DW. The high mass means that the vibration of the lattice cannot be very impressive; hence, the effectiveness of the phonon on the DWR decreases. Also this figure indicates that increasing the sound velocity leads to the reduction of the DWR. The decrease of the resistance is owing to a reduction in the number of excited phonons for high sound velocities in the DW. Therefore, the DWR obtained in the Fe nanowire is higher than that in the Ni nanowire.

## IV. CONCLUSIONS

We have studied the DWR in a magnetic nanowire in the presence of phonons by using the Boltzmann equation, within the relaxation time approximation. The one-dimensional Néel-type DW connecting two regions with antiparallel magnetization was considered. We have investigated how the modulation of the exchange interaction with lattice vibrations and the electron-phonon interaction affect the resistance of the DW. The results indicate that temperature, mass of magnetic ions, and sound velocity in the DW play significant roles. Calculations show that these parameters could have the ability to act on the number of the exited phonons inside the DW; therefore, the DWR is considerably influenced by them. The results indicate that the resistance of the sample enhances as the number of phonons increases. It has also been demonstrated that the exchange interaction strength and deformation potential could significantly affect the resistance of the DWR. Our results clarify that the effect of phonons on the resistance of the DW is not negligible at high temperatures.

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