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Antiferromagnetic and Ferromagnetic Spintronics: Transport in the Two-dimensional Ferromagnet and the Role of In-chain and Inter-chain Interactions

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Antiferromagnetic and ferromagnetic spintronics: transport in the two-dimensional ferromagnet and the role of in-chain and inter-chain interactions

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Abstract. Spin-transport and current-induced torques in ferromagnet heterostructures given by a ferromagnetic domain wall are investigated. Furthermore, the continuum spin conductivity is studied in a frustrated spin system given by the two-dimensional Heisenberg model with ferromagnetic in-chain interaction $J_1 < 0$ between nearest neighbors and antiferromagnetic next-nearest-neighbor in-chain interaction $J_2 > 0$ with aim to investigate the effect of the phase diagram of the critical ion single anisotropy D_c as a function of J_2 on conductivity. We consider the model with the moderate strength of the frustrating parameter such that in-chain spin-spin correlations that are predominantly ferromagnetic. In addition, we consider two inter-chain couplings $J_{\perp,y}$ and $J_{\perp,z}$, corresponding to the two axes perpendicular to chain where ferromagnetic as well as antiferromagnetic interactions are taken into account.

Keywords: spintronics, in-chain interaction, inter-chain interaction

1. Introduction

From the end of 80 decade up to now, the spintronics has witnessed a variety of spin related phenomena such as spin transfer torque, tunneling magnetoresistance[1] and so on. The spintronics demands the spin transport study, where the spin current plays a central role in order of spintronics phenomena to occur in magnetic materials and also in various other materials including semiconductors and oxides.[2, 3] So, the spin current can be used to control the magnetization via spin-transfer-torque and spin-orbit-torque in several magnetic (and also nonmagnetic) materials.[4, 5, 6, 7, 8, 9]

On the other hand, the frustrated antiferromagnetism has been also a rich topic nowadays since to very intriguingly phenomena like topological phase transitions that may occur in this class of systems. The frustration leads to an importance of quantum effects because of the classical order is suppressed and novel phases may occur and to govern the physics at low-energy. Moreover, these systems can exhibit a nematic ground state induced by a spontaneous symmetry breaking induced by terms such as

frustrating interactions in the Heisenberg model.[10] This nematic ordering can occur when spin fluctuations are taken at some axis without any direction being chosen.[11] In a general way, the frustration is present in materials as $LiVCuO_4$, $LiVCu_2O_4$, $LiZrCuO_4$ that are adequately described by ferromagnetic nearest-neighbor interactions J_1 and antiferromagnetic next-nearest-neighbor interactions J_2. [12, 13, 14, 15, 16, 17, 18, 19, 20, 21, 22, 23, 24, 25, 26, 27, 28, 29, 30] All of these materials are of spin-1/2. For spin-1, there are a small number of materials as example the material $NiCl_24SC(NH_2)_2$ which is a quasi-one-dimensional antiferromagnet with easy-plane anisotropy dominating the exchange interaction. [31, 32, 33]

The first generation of spintronic devices are based on spin transport, that utilises the magneto-transport being invented in 2001.[34] The injection efficiency depends on the spin polarisation of the ferromagnet and the spin scattering at the ferromagnetic/non-magnetic interface, where it is also important to eliminate any other effects, namely a stray field from a ferromagnet, that distorts the estimation of the injection efficiency.

The plan of this paper is the following. In section 2, we discuss the spin polarized current. In section 3, we discuss the Meissner mechanism for the spin super current and the antiferromagnetic spintronics for the two-dimensional ferromagnetic model with in-chain and inter-chain interactions. In section 3, we present our results for the spin conductivity. In the last section 4, we present our conclusions and final remarks.

2. Spin polarized current

The spin transport by electrons in a system can be expressed by a spin current in the form $\mathcal{J}_S = -\hbar/2e\,(\mathcal{J}_\uparrow - \mathcal{J}_\downarrow)$, where e is the electron charge. In the same way the charge current $\mathcal{J}_c$ is given by $\mathcal{J}_c = \mathcal{J}_\uparrow + \mathcal{J}_\downarrow$. Both currents obey to a diffusion equation given by

$$\nabla^2(\mu_\uparrow - \mu_\downarrow) = \frac{1}{\Gamma}(\mu_\uparrow - \mu_\downarrow) \tag{1}$$

where $\mu_{\uparrow,\downarrow}$ is the magnetic momentum of each electron and $\Gamma = \sqrt{D\tau}$ is the diffusion coefficient, being D the diffusion constant and τ the spin flip time.[35, 36] Here, we consider the model described by the Hamiltonian

$$\mathcal{H} = -\frac{\hbar^2}{2m}\nabla^2\psi(\mathbf{r}) + \mathbf{V}(\mathbf{r})\psi(\mathbf{r}) + \mathbf{J}\mathbf{s}\cdot\mathbf{S}(\mathbf{r})\psi(\mathbf{r}). \tag{2}$$

where J denotes the exchange integral and $V(\mathbf{r})$ is the potential (nonmagnetic) of the lattice. The last term in the Equation above $J\mathbf{s}\cdot\mathbf{S}(\mathbf{r}')$ provides the interaction electron spin-ferromagnetic domain wall. We consider a homogeneous magnetic domain wall with a collinear magnetization. The interaction among the electrons spins with the spins of the ferromagnetic wall is represented in Fig. 1, where the potential of interaction among the electron spins with the spins of the wall domain is given by $V(\mathbf{r})$. The the aim here is to verify the influence of this on spin wave function of the electrons. The transmission of electron through a domain wall was discussed in Refs. [37, 38]. The purpose here is

to use the Born's expansion for $f_{\uparrow,\downarrow}(\mathbf{k},\mathbf{k}')$ to calculate the effect of interaction electron spin-domain wall on spin wave function of the electrons of the current $\mathcal{J}_c$. The wave function of the electron at large distance from the wall domain is given by

$$|\psi(\mathbf{r})\rangle = \begin{pmatrix} \phi_\uparrow(\mathbf{r}) \\ \phi_\downarrow(\mathbf{r}) \end{pmatrix}, \tag{3}$$

where $\psi_{in}(\mathbf{r}) = \psi(-\infty)$ is the wave function of the electron after the scattering with the wall. Consequently, far from domain wall, we have $\psi_{out}(\mathbf{r})$ given by

$$\begin{aligned} |\psi_\uparrow(\mathbf{k},\mathbf{r})\rangle &= e^{i\mathbf{k}\cdot\mathbf{r}} \begin{pmatrix} 1 \\ 0 \end{pmatrix} + \frac{e^{i\mathbf{k}\cdot\mathbf{r}}}{r} f_\uparrow(\mathbf{k},\mathbf{k}') \begin{pmatrix} 0 \\ 1 \end{pmatrix} \\ |\psi_\downarrow(\mathbf{k},\mathbf{r})\rangle &= e^{i\mathbf{k}\cdot\mathbf{r}} \begin{pmatrix} 0 \\ 1 \end{pmatrix} + \frac{e^{i\mathbf{k}\cdot\mathbf{r}}}{r} f_\downarrow(\mathbf{k},\mathbf{k}') \begin{pmatrix} 1 \\ 0 \end{pmatrix} \end{aligned} \tag{4}$$

where $|\psi_{\uparrow,\downarrow}(\mathbf{k},\mathbf{r})\rangle = \mathbf{S}|\psi_{\uparrow,\downarrow}(\mathbf{k},\mathbf{r})\rangle$ and $\mathbf{S}$ is the scattering matrix.

$$f_{\uparrow,\downarrow}(\mathbf{k},\mathbf{k}') = -\frac{\mu}{2\pi\hbar^2}\langle\psi_{\uparrow,\downarrow}(\mathbf{k},\mathbf{r})|\mathbf{T}|\psi_{\uparrow,\downarrow}(\mathbf{k}',\mathbf{r}')\rangle. \tag{5}$$

$\mathbf{T}$ is given by the Lipmann-Schwinger's equation

$$\mathbf{T} = \mathbf{V} + \mathbf{V}\frac{1}{\omega - \mathcal{H}_0 + i0^+} \tag{6}$$

and

$$\mathcal{H}_0 = -\frac{\hbar^2}{2m}\nabla^2 + V(\mathbf{r}), \tag{7}$$

$$V(\mathbf{r}) = J\mathbf{s}\cdot\mathbf{S}(\mathbf{r}), \tag{8}$$

where we consider $J = 1$. Thus

$$f_{\uparrow,\downarrow}(\mathbf{k},\mathbf{k}') = \sum_{n=0}^{\infty} f^{(n)}_{\uparrow,\downarrow}(\mathbf{k},\mathbf{k}'), \tag{9}$$

where n is the times number that $\mathbf{V}$ enters in the equation above,

$$f^{(1)}_{\uparrow,\downarrow}(\mathbf{k},\mathbf{k}') = -\frac{\mu_B}{2\pi^2\hbar}\langle\psi_{\uparrow,\downarrow}(\mathbf{k},\mathbf{r})|\mathbf{V}|\psi_{\uparrow,\downarrow}(\mathbf{k}',\mathbf{r}')\rangle \tag{10}$$

and the potential $V(x)$ ($\vec{r} = x\hat{i}$) is given by

$$V(x') = -\frac{2J_{sd}}{g\mu_B}\mathbf{s}\cdot\langle\mathbf{S}(x')\rangle, \tag{11}$$

where x' corresponds the region inside of the ferromagnetic wall domain. Consequently, we have

$$f^{(1)}_{\uparrow,\downarrow}(\mathbf{k},\mathbf{k}') = -\frac{2}{e}A(\mathbf{k},\mathbf{k}') \tag{12}$$

where the $A(k)$ coefficient is given by

$$A(k) = \int_{\frac{-\delta}{2}}^{+\frac{\delta}{2}} S\cos(\arctan(e^{-\delta x'}))\sin(kx')dx'. \tag{13}$$

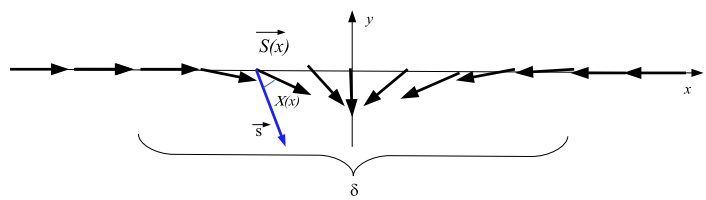

Figure 1. A schematic view of interaction between one electron with the spins of the domain wall of width δ.

The integral above was solved approximately as

$$A(k) \simeq -\frac{1}{4(k^2+\delta^2)}\Big[S\sqrt{2}\Big(e^{\frac{\delta^2}{2}}k\cos\Big(\frac{\delta k}{2}\Big)\delta^3 + e^{\frac{\delta^2}{2}}k^3\cos\Big(\frac{\delta k}{2}\Big)\delta - 4e^{\frac{\delta^2}{2}}k\cos\Big(\frac{\delta k}{2}\Big)\delta$$
$$+e^{\frac{\delta^2}{2}}\sin\Big(\frac{\delta k}{2}\Big)\delta^4 - e^{\frac{\delta^2}{2}}k\sin\Big(\frac{\delta k}{2}\Big)\delta^2k^2 + 2e^{\frac{\delta^2}{2}}\sin\Big(\frac{\delta k}{2}\Big)\delta^2 - 2e^{\frac{\delta^2}{2}}k\sin\Big(\frac{\delta k}{2}\Big)k^2$$
$$+e^{\frac{-\delta^2}{2}}k\cos\Big(\frac{\delta k}{2}\Big)\delta^4 + e^{-\frac{\delta^2}{2}}k^3\cos\Big(\frac{\delta k}{2}\Big)\delta + e^{\frac{-\delta^2}{2}}k\cos\Big(\frac{\delta k}{2}\Big)\delta^4 + e^{-\frac{\delta^2}{2}}k\sin\Big(\frac{\delta k}{2}\Big)\delta^4$$
$$+e^{\frac{-\delta^2}{2}}\sin\Big(\frac{\delta k}{2}\Big)\delta^2k^2 + 2e^{-\frac{\delta^2}{2}}k\sin\Big(\frac{\delta k}{2}\Big)\delta^2 - 2e^{\frac{-\delta^2}{2}}\sin\Big(\frac{\delta k}{2}\Big)k^2\Big)\Big], \tag{14}$$

where δ is the width of the wall. We consider the expansion of the Eq. (9) up to first order. An analysis considering terms of superior order will generate a large quantity of terms in the Eq. (14) and should not generate any change in the scattering. We obtain a very complicated expression for the wave function of the electron after the scattering with the ferromagnetic wall domain however, in a combination of two polarization states. The presence of the coefficient $f(\mathbf{k},\mathbf{k}')$ in the second term making the control of the state of polarization of each electron after the scattering with the domain wall a very difficult problem.

3. Ferromagnetic and antiferromagnetic spintronics

3.1. Meissner effect for the spin supercurrent

It is a fact well known that an example in the nature of local spontaneous breaking of gauge symmetry is the superconductivity. How the charge conductivity is a response to a time-dependent electric field given by Ohm's law, $\mathcal{J}(x,t) = \sigma\vec{E}(x,t)$, in a similar way, we have that the spin current flows in response to a magnetic-field gradient following the Fick's laws for $\nabla\vec{B}(x,t)$ as $\langle\mathcal{J}(x,t)\rangle = \sigma\nabla\vec{B}(x,t)$. Hence, so as for electric superconductors, we should have here $\sigma \to \infty$ for the case of a spin superconductor, where in a finite system with N sites we must have a finite number of spins. As

in general, the spin conductivity cannot be infinite consequently, the gradient of the external magnetic field $\nabla\vec{B}(x,t)$ must be to be zero inside of the a spin superconductor as in the electric superconductor. So, if $\vec{B}$ is zero in the beginning, it must remain zero inside the superconductor even if we apply a gradient of an external magnetic field outside the superconductor. This means that the applied external magnetic field must not dependent on x inside the superconductor. Consequently, the spins of a spin superconductor should generate a current that screens the external gradient of the magnetic field.

The action of the system becomes invariant under the gauge transformation

$$A_\mu(x) \to A_\mu(x) + \partial_\mu\Lambda(x), \tag{15}$$

$$\psi(x) \to \psi(x)e^{iq\Lambda(x)/\hbar}. \tag{16}$$

We introduce the Goldstone boson field $\phi(x)$ that has the property $\phi(x) = \phi(x) + \Lambda(x)$ and

$$\psi(x) = e^{iq\phi(x)/\hbar}\psi(x). \tag{17}$$

where q is the charge. Thus, the magnons in the Heisenberg model must be described by a charged scalar field so as the cooper pairs in the electric superconductor where the Lagrangian describing the interaction between this scalar field with the gradient of the electromagnetic field being given by

$$\mathcal{L} = -\int \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + L_m[A_\mu - \partial_\mu\phi], \tag{18}$$

being L_m a not well known functional.

The Proca's equation is given by

$$\partial_\mu F^{\mu\nu} + \Delta^2 A^\nu = 0 \tag{19}$$

with

$$(\Box + \Delta^2)A^\nu = 0. \tag{20}$$

Thus, the scalar field ψ produces a mass for the photons following the Higg's mechanism where the scalar field that describes the spin waves plays the role of a Higg's field developing a antiferromagnetic vacuum expectation value. So, the photon must acquire a mass Δ inside the spin superconductor like in the electric superconductor with the wave equation becoming the Klein-Gordon equation for the quantity $\mathfrak{B}$ defined by $\mathfrak{B} = \nabla\vec{B}(x,t)$ given by the massive Klein-Gordon equation

$$(\Box + \Delta^2)\mathfrak{B} = 0. \tag{21}$$

Thus, if we apply a gradient of an external magnetic field, the solution of equation above will become $\nabla^2\mathfrak{B} = 0$ for $x < 0$ and $(\nabla^2 - \Delta^2)\mathfrak{B} = 0$ for $x > 0$. The solution for $x > 0$ gives a penetration length given by $l = 1/\Delta$, being the inverse of the mass that the photon acquires inside of the spin superconductor.

3.1.1. Frustrated ferromagnet spin-1/2 Heisenberg chains:

$$\mathcal{H} = J_1 \sum_{\langle i,j\rangle,x} S_i \cdot S_j + J_2 \sum_{[i,j],x} S_i \cdot S_j + J_{\perp,y} \sum_{\langle i,j\rangle,y} S_i \cdot S_j + J_{\perp,z} \sum_{\langle i,j\rangle,y} S_i \cdot S_j, \tag{22}$$

where $\langle i,j\rangle$, x,y,z labels NN bonds along the corresponding axis, and $[i,j],x$ labels NNN bonds along the chain. Furthermore, we consider $J_1 < 0$ and $J_2 \leq 0$, whereas no sign restrictions are valid for $J_{\perp,y}$ and $J_{\perp,z}$.

In the linear response theory, the spin conductivity is the response to an actual frequency-dependent gradient of magnetic field is given by[39, 40, 41]

$$\begin{aligned}
&\langle \mathcal{J}_\beta(\mathbf{k},\omega)\rangle = \sum_\gamma \sigma_{\beta\gamma}(\mathbf{k},\omega) i k_\beta B_\gamma(\mathbf{k},\omega),\\
&\sigma_{\beta\gamma} = \mathrm{Re}(\sigma_{\beta\gamma}) + i\mathrm{Im}(\sigma_{\beta\gamma})\\
&\sigma^{reg}(\omega) = \frac{\mathrm{Im}\{\mathcal{G}(\mathbf{k}=0,\omega)\}}{\omega},
\end{aligned} \tag{23}$$

where we have introduced the spin-flip part of the exchange interaction along the x direction. The Eq. (23) exhibits the desired structure $\langle \mathcal{J}\rangle = \sigma \nabla B^z$, where the formula for the spin conductivity is defined as the linear spin-current response to a uniform, $\mathbf{k} = 0$, frequency-dependent gradient of the magnetic field.[42] In general, the $\mathbf{k} = 0$ conductivity at $T = 0$ may be written as $\mathrm{Re}\left(\sigma_{\beta\gamma}(\omega)\right) = D_S\delta(\omega) + \sigma^{reg}(\omega)$, where D_S is the Drude's weight. Therefore, beyond the $\mathbf{k} = 0$ the relation between the "twist conductivity" and the response to an inhomogeneous magnetic field is not clear.

The Green's function at zero temperature is defined as[40]

$$\mathcal{G}(t) \equiv -\frac{i}{N}\langle 0|\mathbf{T}\mathcal{J}_x(\mathbf{k},t), \mathcal{J}_x(-\mathbf{k},0)|0\rangle. \tag{24}$$

where $\mathbf{T}$ is the time ordering operator. The current-response function $\mathcal{G}(\mathbf{k},\omega)$ at finite temperature is given by

$$\mathcal{G}(\mathbf{k},\omega) = \frac{i}{N}\int_0^\infty dt e^{i\omega t}\langle 0|[\mathcal{J}_x(\mathbf{k},t), \mathcal{J}_x(-\mathbf{k},0)]|0\rangle. \tag{25}$$

$\mathcal{G}(\mathbf{k}=0,\omega\to 0)$ is the susceptibility or retarded Green's function.[40, 41]

The operator for spin current from site j to site $j+x$ is defined by [42, 43, 44, 45, 46]

$$\mathcal{J}_x = i\frac{(J_1 + J_{\perp,y} + J_{\perp,z})}{2}\sum_j (S_j^+ S_{j+x}^- - S_j^- S_{j+x}^+) + \frac{iJ_2}{2}\sum_j (S_j^+ S_{j+2x}^- - S_j^- S_{j+2x}^+) \tag{26}$$

where $j + x$ is the nearest-neighbor site of the site j in the positive x direction. Furthermore, the spin-current operator can be expressed as $\mathcal{J} = -i\left[\mathcal{X}, \mathcal{H}\right]$, where the generating operator $\mathcal{X}$ is $\mathcal{X} \equiv \sum_j x_j S_j^z$, where x_j is the x-coordinate of lattice site j, $\mathcal{J}_x(j) \equiv \mathcal{J}_{x\to j+x}$ and $\mathcal{J}_x \equiv \sum_j \mathcal{J}_x(j)$.

We find the spin current operator in terms of boson operators α and β given by

$$\mathcal{J}_x = t^2 \sum_{\mathbf{k}} \left[\frac{\sin k_x(J_1 + J_{\perp,y} + J_{\perp,z}) + J_2 \sin(2k_x)}{\omega_{\mathbf{k}}}\right](\alpha_{\mathbf{k}} + \beta_{\mathbf{k}})\left(\alpha_{\mathbf{k}}^\dagger + \beta_{\mathbf{k}}^\dagger\right), \tag{27}$$

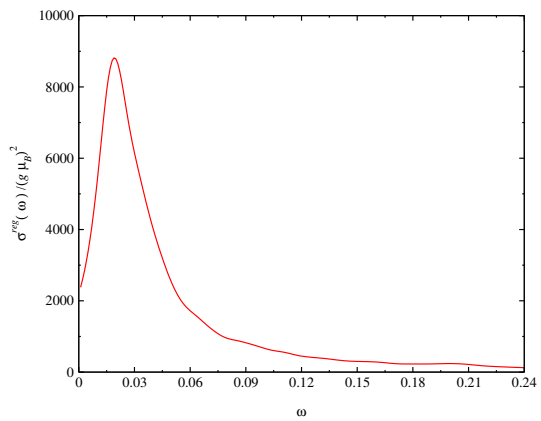

Figure 2. Plot of $\sigma^{reg}(\omega)$ at $T = 0$ using SU(3) Schwinger boson theory). We perform the calculations for values $J_1 = -1.0$, $J_{\perp,y} = 0.1$, $J_{\perp,z} = 0.0$ and $J_2 = 0.2$. The J_2 value is near to the phase transition $0.251 \leq J_{2C} \leq 0.252$. The factor $(g\mu_B)^2$ should be put in case of comparison to experimental data.

where the higher-order terms in Eq. (27) involves terms of four or more boson operators and have been discarded. We find the Green's function given by

$$\mathcal{G}(\mathbf{k},\omega) = t^2(J_1 + J_{\perp,y} + J_{\perp,z})^2 \sum_{\mathbf{k},\mathbf{k}'} \frac{\sin k'_x}{\omega_{\mathbf{k}'}} \frac{\sin k_x}{\omega_{\mathbf{k}}} \mathfrak{g}_{\mathbf{k}\mathbf{k}'}(\omega) + t^2 J_2^2 \sum_{\mathbf{k},\mathbf{k}'} \frac{\sin(2k'_x)}{\omega_{\mathbf{k}'}} \frac{\sin(2k_x)}{\omega_{\mathbf{k}}} \mathfrak{g}_{\mathbf{k}\mathbf{k}'}(\omega) \tag{28}$$

where $\mathfrak{g}_{\mathbf{k}\mathbf{k}'}(\omega)$ is obtained by applying the Wick's theorem following of the Fourier transform after performing a tedious calculation. The Eq. (28) corresponds to the lowest approximation (noninteracting magnetic excitations) which replaces the magnon propagators by the free propagators $\mathcal{G}^0$ and hence, it is valid only in the mean-field approach. Thus, the Green's function for ψ is $\langle \alpha(t)\alpha^\dagger(0)\rangle \to \mathcal{G}_0(t)$ and $\langle \beta^\dagger(t)\beta(0)\rangle \to \tilde{\mathcal{G}}_0(t)$, where $\mathcal{G}_0$ is the bare propagator.

$$\mathcal{G}(\mathbf{k},\omega) = \frac{1}{\pi^2} \int_0^{2\pi} d\omega_1 \mathcal{G}_0(\mathbf{k},\omega_1)\tilde{\mathcal{G}}_0(\mathbf{k},\omega-\omega_1), \tag{29}$$

and

$$\tilde{\mathcal{G}}_0(\mathbf{k},\omega) = -\frac{1}{\omega + \omega_{\mathbf{k}} - i0^+}, \qquad \mathcal{G}_0(\mathbf{k},\omega) = \frac{1}{\omega - \omega_{\mathbf{k}} + i0^+}. \tag{30}$$

Furthermore, we employ the formula[40]

$$\frac{1}{2\pi} \int \mathcal{G}_0(\omega) d\omega \to T \sum_m \mathcal{G}_0(\omega \to i\omega_m), \tag{31}$$

We obtain the retarded Green's function and $\sigma^{reg}(\omega)$, using the SU(3) Schwinger boson theory and applying the Gree-Kubo formula we obtain the continuum conductivity given by

$$\sigma^{reg}(\omega) = t^2 \sum_{\mathbf{k}} \frac{\left[\sin k_x (J_1 + J_{\perp,y} + J_{\perp,z}) + J_2 \sin(2k_x)\right]^2}{\omega_{\mathbf{k}}^3} \delta(\omega - \omega_{\mathbf{k}}), \tag{32}$$

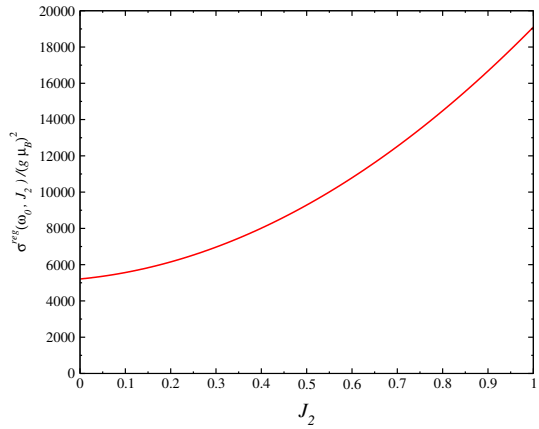

Figure 3. Plot of $\sigma^{reg}(\omega = \omega_0)/(g\mu_B)^2$ as a function of J_2 for $J_1 = -1.0$ and $J_{\perp,y} = 0.1$. We perform the calculations for a ω value near to the peak of the spin conductivity $\omega = \omega_p \approx 0.03$. The factor $(g\mu_B)^2$ should be put in case of comparison to experimental data.

We find the spin conductivity using the SU(3) Schwinger boson theory given as a second-rank tensor, being different from results in literature obtained using the Dyson-Maleev representation which is given by a scalar or zero-order tensor.[40, 41] This difference implies in a different response of the spin current to the gradient of the external magnetic field $\nabla\mathbf{B}$.

We can improve SU(3) Schwinger boson mean-field formalism including the fluctuations around the mean-field result.[47] We can consider the phase fluctuations ϕ_{ij} around the mean-field results for A as $A_{ij} = \tilde{A}e^{i\phi_{ij}}$ taking into account in the action for the Hamiltonian Eq. (.1) which is invariant under the gauge transformation: $b_{ij} \to b_{ij} + (\phi_i - \phi_j)$, $b_j \to b_j e^{i\phi_j}$.

In Fig. 2, we obtain $\sigma^{reg}(\omega)$ at $T = 0$ using the SU(3) Schwinger boson formalism. We perform the calculations for the following values of the parameters: $J_1 = -1.0$, $J_y = 0.1$, $J_z = 0.0$ and $J_2 = 0.2$, where the system is near to the phase transition $0.251 \leq J_{2C} \leq 0.252$. We obtain a large peak for the spin conductivity at $\omega \approx 0.03$ and a finite spin conductivity at $\omega \to 0$ not diverging. In Fig. 3, we obtain $\sigma^{reg}(\omega)$ as a function of J_2 parameter and for the values: $J_1 = -1.0$ and $J_y = 0.1$. We perform the calculations for a ω value near to peak of the spin conductivity $\omega_0 = 0.03$. Due to shape of equation for conductivity, we should obtain the same behavior for the conductivity as a function of J_y. In Fig. 4, we obtain $\sigma^{reg}(\omega)$ as a function of J_2 for $J_1 = -1.0$ and $J_y = 0.1$. We perform the calculations also near to the peak of the spin conductivity, $\omega \approx \omega_0 = 0.03$. In this case, we obtain a monotonically increasing behavior for the conductivity as expected.

4. Conclusions

In brief, we analyze the spin transport and antiferromagnetic and ferromagnetic spintronics. We analyze the case of a lattice model described by the isotropic Heisenberg

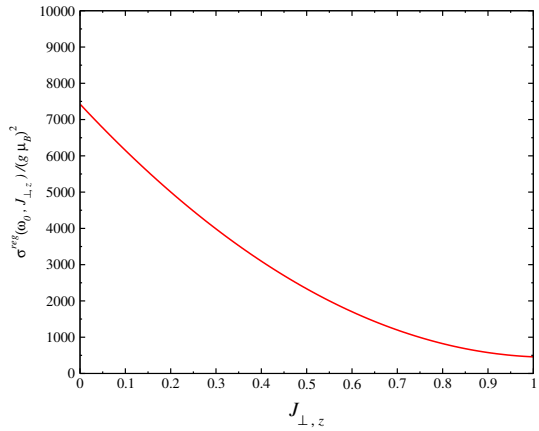

Figure 4. Plot of $\sigma^{reg}(\omega)$ as a function of J_z for $J_1 = -1.0$ and $J_2 = 0.2$. We perform the calculations for a ω value near to the peak of the spin conductivity $\omega = \omega_p \approx 0.03$. The factor $(g\mu_B)^2$ should be put in case of comparison to experimental data.

model with a ferromagnetic in-chain interaction $J_1 < 0$ between nearest neighbors and an antiferromagnetic next-nearest-neighbor in-chain coupling $J_2 > 0$. We obtain a large variation of the spin conductivity with the frustration parameters: J_2, J_x and J_y. There is also a Drude weight for the spin conductivity which exists in many other spin models. However, the purpose here is to analyze the effect of phase transition on continuum conductivity $\omega \neq 0$ where the Drude weight term does not generate any influence. The peak of the spin conductivity can be determined by measuring of magnetization current.[42, 48] In Ref. [49], some experimental techniques were proposed and seem to be feasible. Another experimental technique that can be used is the nuclear magnetic relaxation (NMR). The spin transport in the compound $AgVP_2S_6$ was experimentally investigated using this technique (NMR) in Ref. [50], where the experiment was performed at high temperatures where the behavior for the spin transport is diffusive.

Appendix: SU(3) Schwinger bosons

We consider the model given by

$$\mathcal{H} = J_1 \sum_{\langle i,j\rangle,x} S_i \cdot S_j + J_2 \sum_{[i,j],x} S_i \cdot S_j + J_{\perp,y} \sum_{\langle i,j\rangle,y} S_i \cdot S_j + J_{\perp,z} \sum_{\langle i,j\rangle,y} S_i \cdot S_j + D \sum_i (S_i)^2. \tag{.1}$$

The SU(3) Schwinger boson theory is a theory proposed in Refs. [51, 52] and is well adequate to study the anisotropic XXZ model with single-ion anisotropy D in the range $D > D_c$, for $N \to \infty$ limit where three boson operators t_x, t_y and t_z are defined as

$$t_x^\dagger |v\rangle = |x\rangle, \qquad t_y^\dagger |v\rangle = |y\rangle, \qquad t_z^\dagger |v\rangle = |z\rangle, \tag{.2}$$

where $|v\rangle$ is the vacuum state. In terms of these t_γ operators ($\gamma = x, y, z$), we write the spin-S operators as

$$S^x = -i(t_y^\dagger t_z - t_z^\dagger t_y), \quad S^y = -i(t_z^\dagger t_x - t_x^\dagger t_z), \quad S^z = -i(t_x^\dagger t_y - t_y^\dagger t_x). \tag{.3}$$

We introduces other boson operators $u^\dagger$ and $d^\dagger$ given by

$$u^\dagger = \frac{1}{\sqrt{2}}(t_x^\dagger + it_y), \quad d^\dagger = \frac{1}{\sqrt{2}}(t_x^\dagger - it_y). \tag{.4}$$

Thus

$$|1\rangle = u^\dagger|v\rangle, \quad |0\rangle = t_z^\dagger|v\rangle, \quad |-1\rangle = d^\dagger|v\rangle, \tag{.5}$$

In addition, we impose the constraint condition $u^\dagger u + d^\dagger d + t_z^\dagger t_z = 1$. We employ the further condition that $\langle t_z\rangle = t$.[52] and introduce the Lagrange multiplier $\mu_j(T)$ that is a chemical potential of bosons depending on each site j of the lattice and temperature. From the mean-field approximation, we let $\mu_j(T) = \mu(T)$. The other parameters in theory: t^2, μ, p_1, p_2, p_y and p_z are obtained numerically by a set of integral equations as following

$$2 - t^2 = \frac{1}{N}\sum_{\mathbf{k}} \frac{\Lambda_{\mathbf{k}}}{\omega_{\mathbf{k}}} \coth\left(\frac{\beta\omega_{\mathbf{k}}}{2}\right), \tag{.6}$$

$$\mu = \frac{1}{N}\sum_{\mathbf{k}} \frac{\Lambda_{\mathbf{k}} - \Delta_{\mathbf{k}}}{\omega_{\mathbf{k}}} g_{\mathbf{k}} \coth\left(\frac{\beta\omega_{\mathbf{k}}}{2}\right), \tag{.7}$$

$$p_1 = -\frac{1}{N}\sum_{\mathbf{k}} \frac{\Lambda_k}{\omega_k} \cos k_x \coth\left(\frac{\beta\omega_k}{2}\right), \tag{.8}$$

$$p_2 = -\frac{1}{2N}\sum_{\mathbf{k}} \frac{\Lambda_{\mathbf{k}}}{\omega_{\mathbf{k}}} \cos 2k_x \coth\left(\frac{\beta\omega_{\mathbf{k}}}{2}\right), \tag{.9}$$

$$p_y = -\frac{1}{N}\sum_{\mathbf{k}} \frac{\Lambda_{\mathbf{k}}}{\omega_{\mathbf{k}}} \cos k_y \coth\left(\frac{\beta\omega_{\mathbf{k}}}{2}\right), \tag{.10}$$

$$p_z = -\frac{1}{2N}\sum_{\mathbf{k}} \frac{\Lambda_{\mathbf{k}}}{\omega_{\mathbf{k}}} \cos k_z \coth\left(\frac{\beta\omega_{\mathbf{k}}}{2}\right). \tag{.11}$$

$$\tag{.12}$$

After performing a Fourier transformation followed by the Bogoliubov transformation

$$u_{\mathbf{k}} = \chi_{\mathbf{k}}\alpha_{\mathbf{k}} - \rho_{\mathbf{k}}\beta_{\mathbf{k}}^\dagger, \qquad u_{\mathbf{k}} = \chi_{\mathbf{k}}\beta_{-\mathbf{k}} - \rho_{\mathbf{k}}\alpha_{-\mathbf{k}}^\dagger, \tag{.13}$$

with

$$\chi_{\mathbf{k}} = \sqrt{\frac{\Lambda_{\mathbf{k}} + \omega_{\mathbf{k}}}{2\omega_{\mathbf{k}}}}, \qquad \rho_{\mathbf{k}} = \sqrt{\frac{\Lambda_{\mathbf{k}} - \omega_{\mathbf{k}}}{2\omega_{\mathbf{k}}}} \tag{.14}$$

we find

$$\mathcal{H} = \sum_{\mathbf{k}} \left(\alpha_{\mathbf{k}}^\dagger\alpha_{\mathbf{k}} + \beta_{\mathbf{k}}^\dagger\beta_{\mathbf{k}}\right) + \sum_{\mathbf{k}} (\omega_{\mathbf{k}} - \Lambda_{\mathbf{k}}) + C \tag{.15}$$

with

$$C = \mu N(1-t^2) - \frac{N}{2}(1-t^2)(J_1 + J_2 + J_y + J_z) + 2N\left(J_1p_1^2 + J_2p_2^2 + J_yp_y^2 + J_zp_z^2\right) \tag{.16}$$

$$\Lambda_{\mathbf{k}} = r + 2t^2 g_{\mathbf{k}}, \tag{.17}$$

$$r = (J_1 + J_2 + J_y + J_z)(1-t^2) - \mu + D \tag{.18}$$

$$g_{\mathbf{k}} = J_1\cos k_x + J_2\cos 2k_x + J_y\cos k_y + J_z\cos k_z \tag{.19}$$

$$\Delta_{\mathbf{k}} = 2J_1(t^2 - p_1)\cos k_x + 2J_2(t^2 - p_2)\cos 2k_x + 2J_y(t^2 - p_y)\cos k_y + 2J_z(t^2 - p_z)\cos k_z, \tag{.20}$$

and $\omega_{\mathbf{k}}$ given by

$$\omega_{\mathbf{k}} = \sqrt{\Lambda_{\mathbf{k}}^2 - \Delta_{\mathbf{k}}^2}. \tag{.21}$$

that are the dispersion relation for spin waves. Furthermore, the system present a gap in the spectrum that closes at $\mathbf{k} \in (\pi, \pi)$.

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Figures

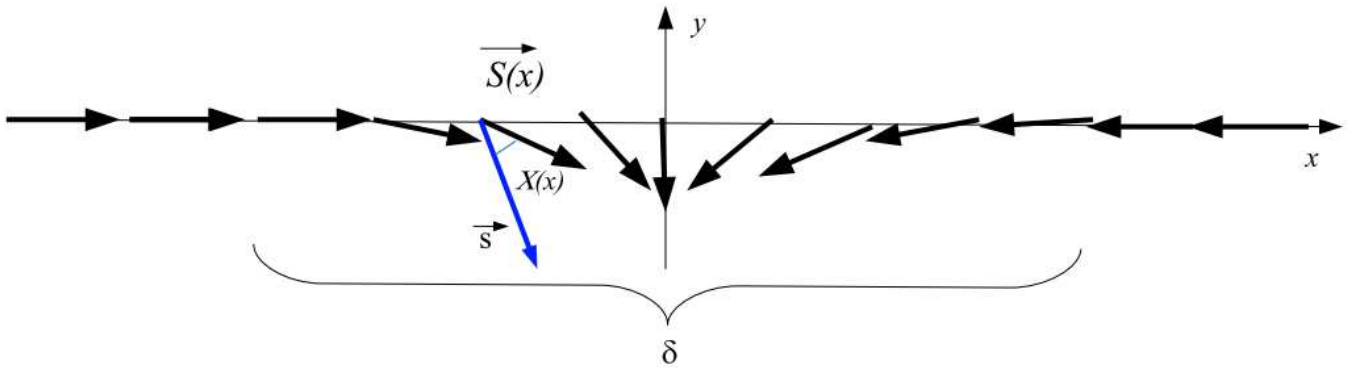

Figure 1

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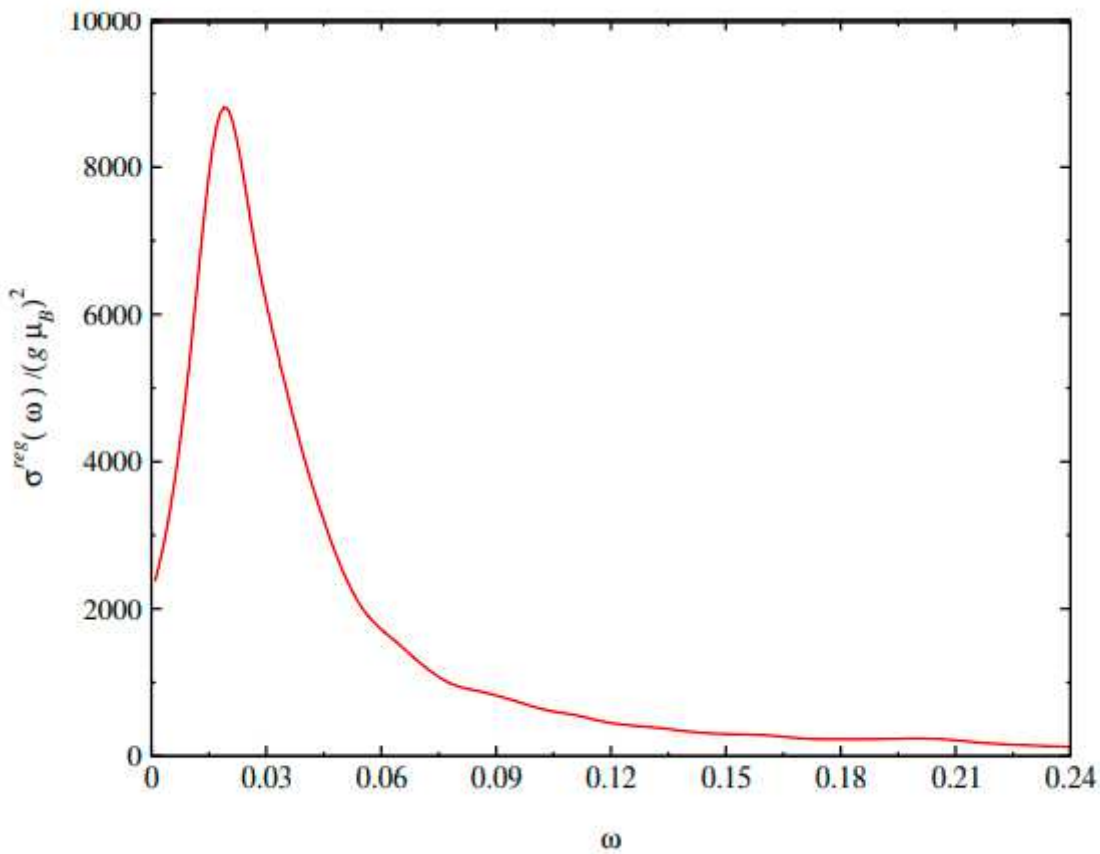

Figure 2

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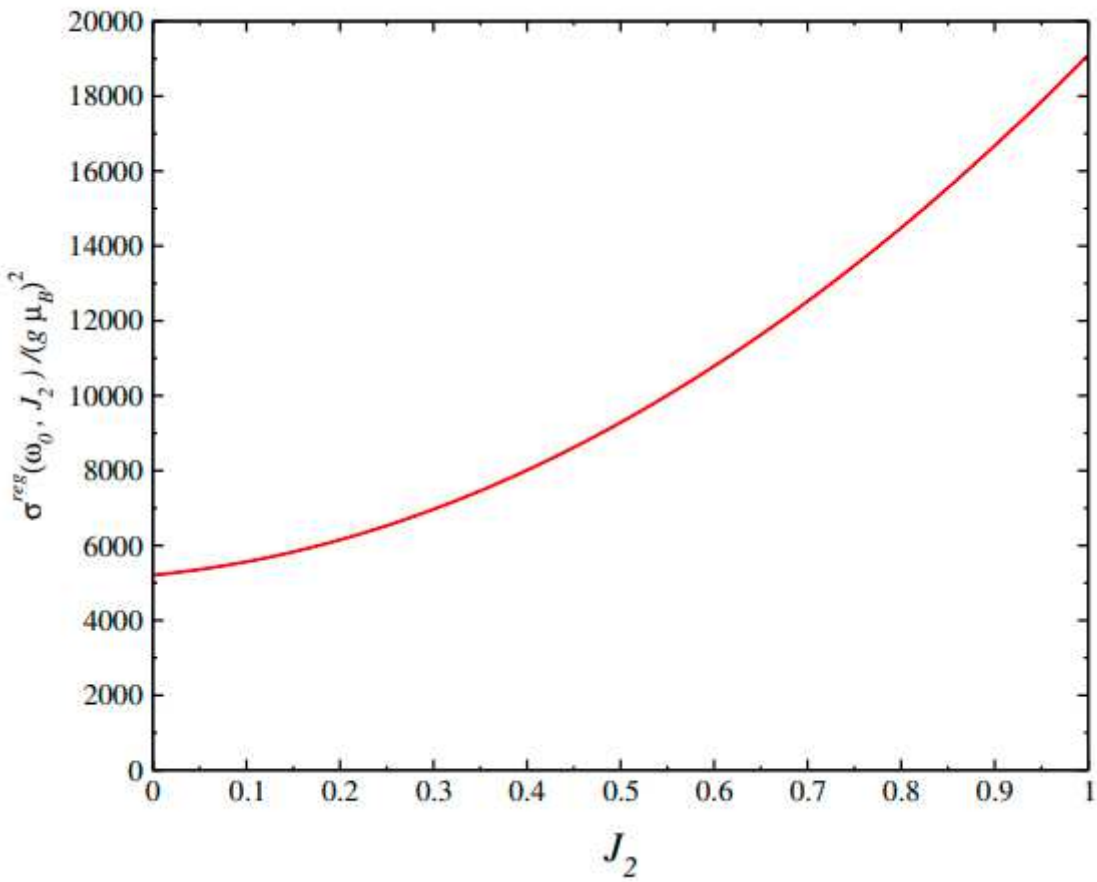

Figure 3

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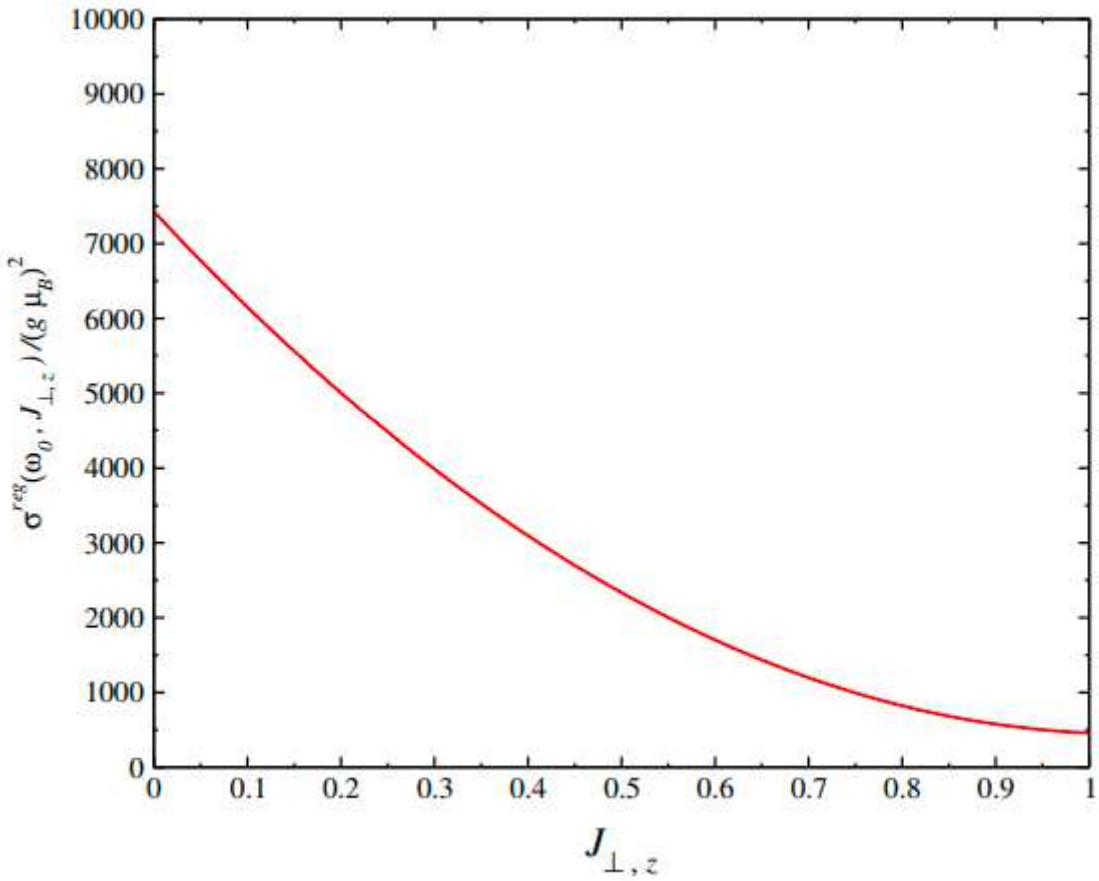

Figure 4

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