Quantum-classical crossover in quasi-one-dimensional systems

C. Bourbonnais
Laboratoire de Physique des Solides, Université de Paris-sud Bâtiment 510, 91405 Orsay, France et Département de Physique, Université de Sherbrooke, Sherbrooke, Québec, Canada J1K 2R1

L. G. Caron
Centre de Recherche en Physique du Solide et Département de Physique, Université de Sherbrooke, Sherbrooke, Québec, Canada J1K 2R1

(Received 25 July 1983; revised manuscript received 6 December 1983)

New results are obtained for quasi-one-dimensional systems. Quantum fluctuations are treated by the Wilson renormalization group and are shown to have a strong influence on the three-dimensional correlations for systems with two-particle interchain coupling. The nonuniversal character of the quantum-classical crossover and the implications on the standard microscopic Ginzburg-Landau theory are examined. We also point out a possible connection to the experimental results on the (tetramethyltetrathiafulvalenium)$_2$-X [(TMTSF)$_2$-X] compounds.

For a large class of quasi-one-dimensional (1D) systems, it is well known that 1D quantum fluctuations can strongly depress the phase-transition temperature T_C. This has clearly been shown[1-3] using the full quantum nature of the 1D correlations[4,5] in a mean-field treatment of the transverse coupling. These calculations neglect, however, the time-dependent effects which are surely relevant whenever the thermal-fluctuation frequencies are less than the characteristic dynamic frequencies.[6] It is the purpose of the present paper to give an adequate description of quantum-fluctuations phenomena in quasi-1D systems. The application to the observed critical behavior of the tetramethyltetrathiafulvalenium-X [(TMTSF)$_2X$] compounds[7] is emphasized.

Consider a set of $N_\perp$ parallel chains of interacting fermions packed into a square lattice with interchain distance $d_\perp$. In the interaction picture, the partition function of the entire system is

$$Z = Z_{||}\left\langle T_\tau \exp\left[-\int_0^\beta H_\perp(\tau)d\tau\right]\right\rangle_{||}, \tag{1}$$

where $\beta = 1/T$, T is the temperature ($k_B = 1$), and τ is the imaginary Matsubara "time." $H_\perp$ is the interchain Hamiltonian. $Z_{||}$ and $\langle\ \rangle_{||}$ are the partition function and the thermodynamic expectation value evaluated with the intrachain Hamiltonian $H_{||}$. By a series of transformations, the single-chain part of the model Hamiltonians that we shall be studying can be written entirely in terms of particle density operators (bosons),[4,5,8]

$$H_{||} = \sum_{i=1}^{N_\perp}\sum_q \left[\frac{2\pi V_F}{L}[\rho_{1,i}(q)\rho_{1,i}(q)+\rho_{2,i}(q)\rho_{2,i}(q)] + \frac{2g}{L}\rho_{1,i}(q)\rho_{2,i}(q)\right]. \tag{2}$$

Here, i is the chain index, L is the length of the chains, V_F is the Fermi velocity of the linearized fermion spectrum, and g is an interaction constant specific to each model ($\hbar = 1$). The operators $\rho_{1(2)}$ refer to the density of particles with velocity V_F ($-V_F$).

These are essentially connected with long-wavelength excitations,[4,5,8] which give power-law singularities in the pair or density-density response function χ_{1D}. At low temperature, for distances x and "times" τ smaller than $\xi_{||,Q} \sim V_\rho/T$ and $1/T$, the correlations are quantum and decay according to a power law

$$\chi_{1D}(x,\tau) \sim |x + iV_\rho\tau|^{\eta-2}.$$

Here, V_ρ is the characteristic velocity of the excitations. At large distances when $x \gg \xi_{||,Q}$ and $\tau \sim 1/T$, we have an exponential decay of correlations $\chi_{1D}(x) \sim T^{2-\eta} e^{-|x|/\xi_{||,T}}$, where $\xi_{||,T} = V_F/T$ is the thermal correlation length. Correspondingly, in Fourier space,

$$\chi_{1D}(q,\omega_m) \sim (\omega_m^2 + V_\rho^2 q^2)^{-\eta/2}$$

for nonvanishing values of $\omega_m = 2\pi m T$, m an integer, and $V_\rho q > 2\pi T$ at $T \neq 0$, whereas $\chi_{1D}(T) \sim T^{-\eta}$ when $m = q = 0$. Since the power-law exponent η and $V_\rho = V_F[1-(g/\pi V_F)^2]^{1/2}$ depend on g, the quantum properties of 1D correlations are nonuniversal.[4,5,8] As $V_\rho \leq V_F$, the thermal (classical) coherence length always encompasses the characteristic quantum length ($\xi_{||,Q} \leq \xi_{||,T}$).

We are interested in a class of interchain coupling which involves two-particle processes. In such cases, $H_\perp$ can be written in the following general form:

$$H_\perp = -\sum_{\alpha=1}^{n}\sum_{i,j}\int V_{ij,\alpha} O_{i,\alpha}^\dagger(x) O_{j,\alpha}(x)dx. \tag{3}$$

$V_{ij,\alpha}$ is the strength of coupling between the chains i and j for the component α of the operators O. For simplicity, V_{ij} will be assumed isotropic. The interchain tunnelings of $2K_F$ electron-hole pairs and of singlet pairs of electrons leading to the charge-density-wave (CDW) transition and to singlet superconductivity (SS) are, respectively, described by the following relevant operators:

$$O_i(x)=\sum_{s=\pm 1}\psi^{\dagger}_{1,i,s}(x)\psi_{2,i,s}(x)$$

and

$$O_i(x)=\sum_{s=\pm 1}s\psi_{1,i,-s}(x)\psi_{2,i,s}(x)\ .$$

Here, the ψ's are the fermion-field operators and s is the spin. Such tunneling for CDW's can be induced by interchain backward scattering,[4,9] or for SS and CDW's from interchain hopping ($t_\perp$) with V proportional to $t_\perp^2$. The latter is only possible for fermions with spin for which a gap $\Delta_s > t_\perp$ exists in the spin excitations.[2-5,10]

A functional integral form for Z can be constructed rigorously via the Hubbard-Stratonovich transformation obtained by the application of the following identity for two commuting operators O and $O^{\dagger}$,[11]

$$e^{O^{\dagger}O}=\int d\psi\, d\psi^{*}\exp[-\pi\,|\,\psi\,|^{2}-\sqrt{\pi}(O^{\dagger}\psi+O\psi^{*})]\ ,$$

to the interaction term of (1). By a cumulant expansion of the remaining thermodynamic average, the result is expressed as a function of the intrachain response functions in all orders of perturbation. These can be approximated following the suggestion of Menyhard.[9,12,13] We arrive at the result $Z=Z_{\|}\int\delta\psi\, e^{-\tilde{H}(\psi,\psi^{*})}$,

$$\tilde{H}(\psi,\psi^{*})=\sum_{\alpha}\sum_{\{\tilde{q}\}}[r_\alpha+A(\omega_m^2+V_\rho^2q^2)^{\eta_\alpha/2}+S_\alpha^2q_\perp^2]\,|\,\psi_\alpha(\tilde{q})\,|^{2}$$
$$+UT/(LN_\perp)\sum_{\alpha\beta}\sum_{\{\tilde{q}\}}\psi_\alpha(\tilde{q}_1)\psi^{*}_\alpha(\tilde{q}_2)\psi_\beta(\tilde{q}_3)\psi^{*}_\beta(\tilde{q}_1-\tilde{q}_2+\tilde{q}_3)+\cdots\ , \tag{4a}$$

where

$$\delta\psi=\prod_{\{\tilde{q}\},\alpha}\frac{d\psi_\alpha(\tilde{q})\,d\psi^{*}_\alpha(\tilde{q})}{\pi\,|\,V_\alpha(\vec{q}_\perp^{\,0})\,|\,\chi_{1D,\alpha}(q,\omega_m)}\ . \tag{4b}$$

Here $(\tilde{q})=(q,\vec{q}_\perp,\omega_m)$ and A is a constant. To this order, $\tilde{H}$ has a quantum Landau-Ginzburg-Wilson (LGW) form. The quadratic term has been obtained by expanding $V_\alpha(\vec{q}_\perp)$ around the transverse ordering wave vector $\vec{q}_\perp^{\,0}$ with S_α as the effective (short) range[14] of V in units of $d_\perp$, and also by the above-mentioned asymptotic behavior of χ_{1D} for the 1D quantum-fluctuation phase space

$$\omega_D\geq\begin{vmatrix}\omega_m\\ V_\rho q\end{vmatrix}>2\pi T$$

(see Refs. 4, 5, 8, 10, and 12). Here, $r_\alpha=(T/T_{C,\alpha}^{MF})^{\eta_\alpha}-1$, with $T_{C,\alpha}^{MF}\sim|\,V(q_\perp^0)\,|^{1/\eta_\alpha}$ as the three-dimensional (3D) mean-field (MF) transition temperature which is determined by the condition

$$\chi_{1D}^{-1}(T_{C,\alpha}^{MF})\,|\,V_\alpha(\vec{q}_\perp^{\,0})\,|^{-1}-1=0\ .$$

The mode-mode coupling strength U is proportional to the fourth-order 1D free-fermion loop at $\omega=q=0$, which is finite at $T\neq 0$.[9,12] The natural energy cutoff for quantum effects $\omega_D=V_\rho q_0$ is the Debye energy of the 1D collective excitations,[4,5] with q_0 as the cutoff wave vector. In the transverse direction $q_{\perp 0}=2\pi/d_\perp$ is the natural cutoff. Clearly, the above quasi-1D LGW form for the ψ field has an effective anisotropic range of interaction which can be long range in space and time for the longitudinal direction ($\eta<2$) and short range in the transverse directions.[15]

In the cases where $T_C^{MF}\ll\omega_D$, the static approximation for $\tilde{H}$ is not justified.[6,16] It must be preceded by the integration of all quantum degrees of freedom. This can be done properly by applying the Wilson renormalization-group (RG) technique for quantum functionals (we shall denote these techniques as QRG),[6(b),16] first introduced by Béal-Monod. Following Ref. 6(b), we first integrate the ψ's of the outer-shell phase space

$$e^{-2l}<\delta(\omega^2+q^2)^{\eta_\alpha/2}+S_\alpha^2q_\perp^2<1$$

in reduced units ($\delta>1$), and we rescale anisotropically the variables $q\Rightarrow q'=e^{2l/\eta_\alpha}q$, $\vec{q}_\perp\Rightarrow\vec{q}\,'_\perp=e^{l}\vec{q}_\perp$, and $\omega\Rightarrow\omega'=e^{2l/\eta_\alpha}\omega$ taken to be continuous. The fields are then rescaled as

$$\psi_\alpha\Rightarrow\psi'_\alpha=\psi_\alpha(\tilde{q})e^{-(2/\eta_\alpha+d/2+1/2)l}$$

in order to preserve the form of $\tilde{H}$, the phase space, and field densities.[14] Here, d is the dimension and l is the infinitesimal generator of the RG. This rescaling procedure leads to the usual forms of the RG equations,[6,14]

$$\frac{dr_\alpha}{dl}=2r_\alpha+\frac{K_{d,n}U}{1+r_\alpha}\ , \tag{5a}$$

$$\frac{dU}{dl}=\epsilon_\alpha U-\frac{K'_{d,n}U^2}{(1+r_\alpha)^2}\ . \tag{5b}$$

Here, $K_{d,n}$ and $K'_{d,n}$ are constants which depend on the dimension d and the number n of the ψ components. The rescaling gives a different value of $\epsilon_\alpha=4-(d-1+4/\eta_\alpha)$ for $\eta_\alpha<2$. In contrast to the classical case,[14] where $\eta=2$, $\omega=0$, and $\epsilon=4-d$, the effective dimensionality in the critical-behavior sense is increased by anisotropic quantum effects to $d_\alpha^{*}=d-1+4/\eta_\alpha>4$ for $\eta_\alpha<2$ and $d=3$. If $\epsilon_\alpha<0$, the quartic term in H acts as a small perturbation for the relevant 3D Gaussian properties.[14] Higher-order terms in H do not affect this behavior since they have a faster decay during the renormalization, namely $U'_K=U_Ke^{\epsilon_\alpha^K l}$ with

$$\epsilon_\alpha^K=4-2(K/\eta_\alpha-2/\eta_\alpha+1)<\epsilon_\alpha^{K-2}$$

for $K\geq 6$ and $d=3$. The passage to a classical thermodynamic regime can be determined when only one finite Matsubara frequency remains,[6(b)] that is, when $\omega_De^{-2l_\alpha^{*}/\eta_\alpha}=2\pi T$. This defines two characteristic wave vectors in the fluctuation phase space,

$$q_{c,\alpha}=q_0 e^{-2l^*_\alpha/\eta_\alpha}=(2\pi T/\omega_D)q_0$$

and

$$q_{\perp c,\alpha}=(2\pi T/\omega_D)^{\eta_\alpha/2}q_{\perp 0}\ .$$

For $\vec{q}>\vec{q}_{\perp c,\alpha}$ and $q>q_{\perp c,\alpha}$, the fluctuations are essentially quantum and they are governed by the QRG equations (5). This indicates, in contrast to the usual static calculation, the possiblity of transverse propagation of the quantum correlations. On the other hand, for $q<q_{c,\alpha}$ and $\vec{q}_\perp<\vec{q}_{\perp c,\alpha}$, the fluctuations are static and classical with $\epsilon=4-d$.

In the classical regime, the critical behavior is treated by the usual classical LGW functional with the renormalized values of T_c and U and the new cutoff q_c and $\vec{q}_{\perp c}$.[6,14] To first order in U, the renormalization of T_c is given by the solution of (5),

$$(T^{\rm MF}_{c,\alpha}-\bar{T}_{c,\alpha})/T^{\rm MF}_{c,\alpha}=1-(|\bar{V}_\alpha|/|V_\alpha|)^{1/\eta_\alpha}\ , \tag{6}$$

with

$$|\bar{V}_\alpha|=|V_\alpha|\left[1-\left[\frac{UK_{d,m}}{|\epsilon_\alpha|+2}\right](1-e^{-(|\epsilon_\alpha|+2)l^*_\alpha})\right]\ , \tag{7}$$

where

$$l^*_\alpha=\tfrac{1}{2}\eta_\alpha\ln(\omega_D/2\pi T^*_\alpha)\ .$$

The quantum-classical crossover temperature T^*_α is reached when the transverse coherence length $\xi_{\perp,\alpha}$ is of the order of the transverse characteristic quantum length $d_{\perp c,\alpha}\equiv 2\pi q^{-1}_{\perp c,\alpha}$. For smaller values of $\xi_{\perp,\alpha}$ the correlations are governed by quantum dynamics, while for larger values they are governed by thermal fluctuations. Since for $T>T^*_\alpha$, $\epsilon_\alpha<0$, a Gaussian-like behavior of $\xi_{\perp,\alpha}$ leads to $\xi_{\perp,\alpha}=\xi_{\perp 0,\alpha}(r_\alpha)^{-1/2}$ with $\xi_{\perp 0,\alpha}=S_\alpha/\sqrt{2}d_1$. The crossover condition $\xi_{\perp,\alpha}(T^*_\alpha)=d_{\perp c,\alpha}$ yields, in the linear approximation for r_α,

$$\Delta t^*_\alpha\equiv(T^*_\alpha-\bar{T}_{c,\alpha})/\bar{T}_{c,\alpha}$$
$$\simeq\tfrac{1}{2}(S^2_\alpha/\eta_\alpha)(2\pi\bar{T}_{c,\alpha}/\omega_D)^{\eta_\alpha}/[1-(S^2_\alpha/2)(2\pi\bar{T}_{c,\alpha}/\omega_D)^{\eta_\alpha}]. \tag{8}$$

The crossover exponent becomes $\phi_\alpha=1/\eta_\alpha$, with the temperature as the symmetry-breaking parameter.[6(b)] Δt^*_α is then a measure of the range of validity of the classical LGW treatment. The region $T<\bar{T}_{c,\alpha}$ is subjected to the same conditions since quantum effects for $q_\perp>q_{\perp c,\alpha}$ persist throughout the long-range-order regime. As η_α depends on g, ϕ_α reflects the *nonuniversal* way that the system loses its Gaussian quantum properties. From the crossover condition on the longitudinal coherence length, $\xi_{||,\alpha}(T=T^*_\alpha)=2\pi q^{-1}_{c,\alpha}\equiv d_{||c,\alpha}$, and Eq. (8), it follows that

$$\xi_{||,\alpha}=\xi_{||0,\alpha}[(T-\bar{T}_{c,\alpha})/\bar{T}_{c,\alpha}]^{-1/\eta_\alpha}$$

for $T\gtrsim T^*_\alpha$, which might have been guessed in advance from (4).[15] More generally, for arbitrary interaction in a transverse direction[15] $q^2_\perp\to q^\sigma_\perp$ in (4) $(0<\sigma\leq 2)$, $\phi_\alpha=1/\eta_\alpha$, and the longitudinal and the transverse exponents $\nu_{||,\alpha}=1/\eta_\alpha$ and $\nu_{\perp,\alpha}=1/\sigma_\alpha$ for ξ satisfy the following anisotropic relation: $\nu_{\perp,\alpha}/\nu_{||,\alpha}=\eta_\alpha/\sigma_\alpha$ for the growth of the quasi-1D quantum-correlation cluster at $T>T^*$.

A few remarks are now in order. From Eqs. (5)—(8), in the 1D free-electron-gas limit where $\eta_\alpha\to 0$, $|\epsilon_\alpha|\to\infty$, and $l^*_\alpha\to 0$, the correction to $T^{\rm MF}_{C,\alpha}$ vanishes and $T^*_\alpha\gg T^{\rm MF}_{c,\alpha}$. Similarly, near the classical limit, when $\omega_D\lesssim 2\pi T^{\rm MF}_{c,\alpha}$ or $\eta_\alpha\approx 2$, we also have $\bar{T}_c\to T^{\rm MF}_{c,\alpha}$. For both cases, an unrenormalized static LGW treatment becomes justified. Furthermore, one must note that the limit $\eta\to 0$ with $d^*\to\infty$ is consistent with static BCS-like thermodynamics near T_c.[12] However, for other values of η and ω_D it is possible that the renormalization of T_c due to the addition of a large number of nonthermal-fluctuating modes may be of the order of $T^{\rm MF}_{c,\alpha}$ itself. This can occur despite the fact that in the renormalized static calculation the relevance of U in the Ginzburg sense is restricted to a narrow temperature range around $T^{\rm MF}_{c,\alpha}$.[9,12] Similar large quantum-fluctuation corrections to LGW parameters have been already obtained in the context of itinerant ferromagnetism.[17] Clearly, in such cases, the use of the mean-field transition temperature is surely not valid, and even its low-order corrections can be questionable, and it is more preferable to consider $\bar{T}_{c,\alpha}$ as a parameter of the problem.

We now look at some phase transitions where the results derived above are applicable (i) A backward scattering process for V_{ij} can give rise to a CDW transition.[5,9] In the case of spinless fermions the short-wavelength cutoff $\Lambda_0=2\pi q_0^{-1}$ is the range of the interaction which can be taken to be of the order of the 1D lattice constant a.[8] For weak coupling,[4,8] $V_\rho\simeq V_F$ and $0<\eta\leq 1$. Quantum corrections increase with η. Since there is no gap in the 1D interacting Luttinger model,[8] the ellipsoid-shaped volume $V_c\sim\pi d^2_{\perp c}d_{||c}$ spanned by quantum fluctuations can be considered to be the real size of the 3D electron-hole pairs.[18] Their collective motion which drives the transition is then described by the classical LGW treatment. (i) For the SS transition and weak attractive coupling along the chains, $0<\eta\leq 1$ and $V_\rho\simeq V_F$, and Λ_0 is increased to $\xi_0=V_F/\Delta_s$, which is of the order of the size of the bound pairs of electrons.[5] With 1D strong local attractive coupling between electrons, quantum effects can be stronger since $\xi_0\simeq a$, $0<\eta\lesssim\frac{3}{2}$, and $V_F>V_\rho\gtrsim\frac{1}{4}V_F$. With a nonretarded attractive coupling between electrons and $V_\alpha\sim t^2_\perp$, some competitive effects between the SS and CDW types of long-range ordering can occur. According to Refs. 4 and 5 it is the strongest in the weak attractive coupling where

$$0<\eta_{\rm CDW}<\eta_{\rm SS}\ \text{or}\ T^{\rm MF}_{c,\rm CDW}<T^{\rm MF}_{c,\rm SS}\ ,$$

whereas it is small in the strong coupling case where[5] $\eta_{\rm CDW}=0$ when $\eta_{\rm SS}=\frac{3}{2}$. Thus, we have the possibility of a quantum-classical crossover for both types of correlation with $\phi_{\rm CDW}=1/\eta_{\rm CDW}$, $\phi_{\rm SS}=1/\eta_{\rm SS}$, and $\Delta t^*_{\rm CDW}\geq\Delta t^*_{\rm SS}$. The classical regime can be described by retaining only the correlations with the highest $\bar{T}_c$.[19] (iii) When the 1D electron-phonon interaction dominates the 1D Coulomb interaction g, a Peierls transition can occur with $V\sim t^2_\perp$ (or transverse backward scattering), $V_\rho=(m/m^*)^{1/2}V_F$, $\Lambda_0=V_F/\Delta_0$, and $d_{\perp c}=T^0_c/T^*(m/m^*)^{1/2}d_\perp$. m^* is the

CDW effective mass and m is the electron-band mass. In the adiabatic approximation $m^* \gg m$ ($\eta \approx 2$),[20] the gap $\Delta_0 = 1.73 T_{c0}$, where T_{c0} is the 1D mean-field transition temperature. It follows that the quantum correlations to the transition are quite small, namely $0 < \omega_D \lesssim T_c^{\rm MF}$, $l^* \approx 0$, and $d_{\perp c} \approx d_\perp$. This means that the use of an unrenormalized classical LGW functional becomes justified for essentially all temperatures. The conclusions are completely different if we are in the *nonadiabatic* limit when $m^* \approx m$. In this case, the electron-phonon interaction is equivalent to a nonretarded attractive coupling between electrons, and therefore the treatment in (ii) applies. However, the above discussion was restricted to the incommensurate cases and it appears that the conclusions for the *half-filled band,* or *commensurate,* case are different. For example, in the nonadiabatic limit of electron-phonon interaction, because of the relevance of umklapp scattering processes,[21] the 1D problem is highly quantum[4] and there is a gap in both charge and spin degrees of freedom with only divergent CDW correlations. These are characterized[4] by $\eta = 2$, but with $\omega_D \gg T_c^{\rm MF}$. With a 3D kinetic or potential coupling it follows that $\epsilon = 0$, $d^* = 4$ in the quantum regime ($T^* < T$), and $\phi = \frac{1}{2}$, for $\Delta t^* \ll 1$; that is, strong quantum corrections must be taken into account.

The implications of the above results for the $(TMTSF)_2X$ family of superconductors[7] are of interest. This is especially true if one looks at the effect of the renormalization on the Ginzburg[22] critical width in units of T_c,

$$\Delta t_G = \Delta t_G^0 (d_\perp / d_{\perp c})^4 (\bar{T}_C / T_C^{\rm MF})^2 \propto q_{\perp c}^4 q_c^2 \bar{T}_C^2 U^2 ,$$

where Δt_G^0 comes from the unrenormalized calculation, which have been discussed in detail in Ref. 12. Several arguments suggest that for organic superconductors,[4,7,12,23] $1 > \eta > 0$ and $\Delta_s = 20$–30 K. From our analysis, $d_{\perp c}/d_\perp = (\Delta_s / T^*)^{\eta/2}$ and $\Delta t^* \simeq S^2/2$ ($\bar{T}_C / \Delta_s$), and from the experimental value of $\bar{T}_C \approx 1$ K we can conclude that the quantum corrections to the static prediction ($\eta = 0$) $0 \lesssim \Delta t_G^0 \lesssim 0.05$ ($\eta \approx 1$) can be important; that is, $\Delta t_G \ll \Delta t_G^0$ and the chances of observing a real departure from 3D mean-field behavior are small. The electronic specific heat is linear in temperature when it is dominated by 1D quantum-fluctuation effects[4,5,10,12,24] for $T_C < T < \Delta_s$. Therefore, the transition would resemble that of an ordinary 3D superconductor despite the quasiparticle confinement due to a pseudogap of width $2\Delta_s$.[12,25] The small width of the transition observed by specific-heat measurements on $(TMTSF)_2ClO_4$ (Ref. 26) is therefore compatible with a quantum-fluctuation picture.[7]

As a final remark, we wish to emphasize that the present approach does not apply when $\Delta_s \ll t_\perp$. When spin- or charge-excitation gaps are not relevant in the presence of transverse hopping, transverse *one-particle* events must be taken into account. As stated previously, this was neglected by writing $H_\perp$ in the form given by (3).

In conclusion, we have pointed out that 1D quantum fluctuations can propagate in the transverse direction. For a large class of quasi-1D Hamiltonians, a RG treatment for these fluctuations leads to a nonuniversal crossover from a boson variable description, valid when 1D quantum effects dominate, to a standard, but renormalized, static LGW approach near the transition.

The authors would like to thank M. T. Béal-Monod, H. J. Schulz, D. Jérome, J. Friedel, P. Pfeuty, and J. Solyom for several discussions and comments. We also thank S. Barisic and K. B. Efetov for a number of enlightening discussions about some intrachain aspects of the problem. This work was supported by the Natural Sciences and Engineering Council of Canada and Le Ministère de l'Education du Québec.

[1] Y. Imry, P. Pincus, and D. J. Scalapino, Phys. Rev. B 12, 1978 (1975).

[2] V. J. Emery, Phys. Rev. B 14, 2989 (1976).

[3] L. P. Gorkov and I. E. Dzyaloshinskii, Zh. Eksp. Teor. Fiz. 67, 397 (1974) [Sov. Phys.—JETP 40, 198 (1975)]; V. N. Prigodin and Y. A. Firsov, *ibid.* 76, 1602 (1979) [49, 813 (1979)].

[4] J. Solyom, Adv. Phys. 28, 201 (1979); V. J. Emery, in *Highly Conducting One-Dimensional Solids,* edited by J. T. Devreese, R. P. Evrard, and V. E. van Doren (Plenum, New York, 1979), p. 247; D. Jerome and H. J. Schulz, Adv. Phys. 31, 299 (1982).

[5] K. B. Efetov and A. I. Larkin, Zh. Eksp. Teor. Fiz. 69, 764 (1975) [Sov. Phys.—JETP 42, 390 (1975)].

[6] (a) L. D. Landau and E. M. Lifshitz, *Statistical Physics* (Pergamon, Oxford, 1958), Sec. 112; (b) J. A. Hertz, Phys. Rev. B 14, 1165 (1976).

[7] (a) D. Jerome, A. Mazaud, M. Ribault, and K. Bechgaard, J. Phys. (Paris) Lett. 41, L95 (1980); (b) Proceedings of the International Colloquium on the Physics and Chemistry of Synthetic and Organic Metals, Les Arcs, France [J. Phys. (Paris) Colloq. 44, C3-775 (1983)].

[8] A. Luther and I. Peschel, Phys. Rev. B 9, 2911 (1974); 12, 3908 (1975).

[9] K. A. Klemm and H. Gutfreund, Phys. Rev. B 14, 1086 (1976); N. Menyhard, J. Phys. C 11, 2207 (1978).

[10] K. B. Efetov, Zh. Eksp. Teor. Fiz. 70, 2339 (1976) [Sov. Phys.—JETP 43, 1221 (1967)].

[11] J. Hubbard, Phys. Rev. Lett. 3, 77 (1959); R. L. Stratonovich, Dokl. Akad. Nauk SSSR 115, 1097 (1957) [Sov. Phys.—Dokl. 2, 416 (1957)]; D. J. Amit and H. Keiter, J. Low Temp. Phys. 11, 603 (1973).

[12] H. J. Schulz and C. Bourbonnais, Phys. Rev. B 27, 5856 (1983).

[13] A. A. Abrikosov, L. P. Gorkov, and I. E. Dzyaloshinskii, *Method of Quantum Field Theory in Statistical Physics* (Prentice-Hall, Englewood Cliffs, N.J., 1963).

[14] K. G. Wilson and J. Kogut, Phys. Rep. 12, 75 (1974); M. E. Fisher, Rev. Mod. Phys. 46, 597 (1974).

[15] M. E. Fisher, S. k. Ma, and B. G. Nickel, Phys. Rev. Lett. 29, 917 (1972).

[16] M. T. Béal-Monod, Solid State Commun. 14, 667 (1974); M. T. Béal-Monod and K. Maki, Phys. Rev. Lett. 34, 461 (1975).

[17] K. K. Murata, Phys. Rev. B 12, 282 (1975).

[18] C. Bourbonnais and L. G. Caron, in Proceedings of the International Colloquium on the Physics and Chemistry of Synthetic and Organic Metals, Les Arcs, France [J. Phys. (Paris)

Colloq. 44, C3-911 (1983)].

[19]For the special case $\eta_{CDW}=\eta_{SS}$ with $V_{CDW}=V_{SS}\sim t_\perp^2$, a very small additional 3D coupling will reduce the symmetry of the order parameter in the classical regime. For example, if we add small transverse backscattering to V_{CDW}, the reduction of symmetry CDW + SS→CDW can be completely described by the classical crossover phenomena of the anisotropic n-vector model. See, for example, M. E. Fisher and P. Pfeuty, Phys. Rev. B 6, 1889 (1972). The results of (ii) can be transposed directly to the Heisenberg-Ising antiferromagnetic transition with the following relations, SS $\Longleftrightarrow XY$ and CDQ $\Longleftrightarrow Z$, between correlations (see Refs. 2 and 10).

[20]H. Fukuyama, J. Phys. Soc. Jpn. 41, 513 (1976).

[21]L. G. Caron and C. Bourbonnais, Phys. Rev. B 29, 4230 (1984).

[22]G. L. Ginzburg, Fiz. Tverd. Tela (Leningrad) 2, 2031 (1960) [Sov. Phys.—Solid State 2, 1824 (1961)].

[23]K. B. Efetov, J. Phys. (Paris) Lett. 44, L369 (1983).

[24]H. Fukuyama, T. M. Rice, C. M. Varma, and B. I. Halperin, Phys. Rev. B 10, 3775 (1974).

[25]Y. Suzumura, Prog. Theor. Phys. 63, 41 (1980).

[26]P. Garoche, R. Brusetti, and D. Jerome, J. Phys. (Paris) Lett. 43, L147 (1982).