Magnetization distribution in the transverse Ising chain with energy flux

V. Eisler,1 Z. Rácz,1,2 and F. van Wijland2

1Institute for Theoretical Physics, Eötvös University, Pázmány sétány 1/a, 1117 Budapest, Hungary
2Laboratoire de Physique Théorique, Bâtiment 210, Université de Paris–Sud, 91405 Orsay Cedex, France

(Received 28 January 2003; published 27 May 2003)

The zero-temperature transverse Ising chain carrying an energy flux \( j_E \) is studied with the aim of determining the nonequilibrium distribution functions, \( P(M_x) \) and \( P(M_y) \) of its transverse and longitudinal magnetizations, respectively. An exact calculation reveals that \( P(M_y) \) is a Gaussian both at \( j_E = 0 \) and at \( j_E \neq 0 \), and the width of the distribution decreases with increasing energy flux. The distribution of the order-parameter fluctuations, \( P(M_x) \), is evaluated numerically for spin chains of up to 20 spins. For the equilibrium case \( (j_E = 0) \), we find the expected Gaussian fluctuations away from the critical point, while the critical order-parameter fluctuations are shown to be non-Gaussian with a scaling function \( \Phi(x) = \Phi(M_x/\langle M_y \rangle) = \langle M_y \rangle P(M_x) \) strongly dependent on the boundary conditions. When \( j_E \neq 0 \), the system displays long-range, oscillating correlations but \( P(M_x) \) is a Gaussian nevertheless, and the width of the Gaussian decreases with increasing \( j_E \). In particular, we find that, at critical transverse field, the width has a \( j_E^{-3/8} \) asymptotic in the \( j_E \rightarrow 0 \) limit.

DOI: 10.1103/PhysRevE.67.056129 PACS number(s): 05.50.+q, 05.60.Gg, 05.70.Ln, 75.10.Jm

I. INTRODUCTION

Nonequilibrium steady states (NESS) have been much studied but a description of some generality has not emerged so far. Among the many approaches tried, there is one that continues to receive particular attention. It is an attempt to understand the general features of NESS through studies of nonequilibrium phase transitions [1–3]. The basic assumption here is that the universality displayed by equilibrium phase transitions carries over to critical phenomena in NESS, as well. Thus, by investigating the similarities and differences from equilibrium, one may gain an understanding of the role of various components of the competing dynamics generating the steady state. For example, one may find by observing the universality classes of various nonequilibrium phase transitions that dynamical anisotropies often yield dipolelike effective interactions [4–6] or that competing nonlocal dynamics (anomalous diffusion) generates long-range, power-law interactions [7].

The extension of concepts of critical phenomena to phase transitions in NESS also implies that the distribution functions of macroscopic quantities (such as the order parameter) are nontrivial (non-Gaussian). They are universal, however, and characterize the given nonequilibrium universality class, too. The advantage of studying the scaling functions associated with the distribution functions is that building these functions does not involve any fitting procedure and thus they allow for a fit-free comparison with experiments. Indeed, using these distribution functions, a number of interesting results have been derived for surface growth as well as for other nonequilibrium processes [8–13].

In this paper, we continue our studies of nonequilibrium distribution functions by investigating the effects of a nonequilibrium constraint on a well-known quantum phase transition; namely, we take the transverse Ising chain that has an order-disorder transition as the transverse field \( h \) is varied, and drive it by a field to produce an energy flux \( j_E \) through the system. The resulting steady states have been described in Ref. [14] and it has been found that, in addition to the equilibrium phases, a flux-carrying nonequilibrium phase appears, which is distinct by its correlations decaying with distance as a power law. We shall be concerned with the distribution function in the various phases of the above system. More precisely, we shall determine the steady-state distribution functions \( P(M_x) \) and \( P(M_y) \) of the \( M_x \) (nonordering field) and \( M_y \) (ordering field) components of the macroscopic magnetization in all three phases of the system at and near its critical point.

The results are surprisingly simple. The distribution functions are Gaussian in the equilibrium phases away from the critical point. This is expected since we have macroscopic quantity and the correlations decay exponentially. The distribution of the nonordering field remains a Gaussian at the critical point of the equilibrium system, as well. The reason for this is that although the appropriate correlations decay with distance \( n \) as a power law but the exponent in the power is large \((1/n^2)\), so that the fluctuations \( \langle M_x^2 \rangle - \langle M_x \rangle^2 \) do not diverge at \( h = h_c \). The distribution function of the ordering field becomes nontrivial at \( h_c \) and our numerical calculations demonstrate that \( P(M_y) \) depends strongly on the boundary conditions taken to be periodic, antiperiodic, and free. The unexpected simplicity is in the current-carrying phase where the energy flux generates long-range correlations decaying as a power law \((1/\sqrt{n})\) but, nevertheless, the distributions are Gaussian. The mathematical reason for this lies in the oscillating character of the correlations, which prevents the divergence of the spatial sum of the correlations which in turn are proportional to the fluctuations. Physically, the oscillations in the correlations can be traced to the form of energy flux [see Eq. (3) below], which suggest that the consecutive \( x \) and \( y \) components of the spins are more and more rigidly interconnected as \( j_E \) is increased and thus fluctuations decrease with increasing \( j_E \). This picture will be seen to be valid near the nonequilibrium phase boundaries where the fluctuations as a function of \( j_E \) can be explicitly calculated.

1063-651X/2003/67(5)/056129(8)/$20.00 67 056129-1 ©2003 The American Physical Society
II. TRANSVERSE ISING MODEL WITH ENERGY FLUX

The transverse Ising chain is one of the simplest systems displaying a critical order-disorder transition [14–19]. It is defined by the Hamiltonian

$$\hat{H} = -J \sum_i s_i^x s_{i+1}^x - \frac{h}{2} \sum_i s_i^z,$$

(1)

where $s_i = \frac{1}{2} \sigma_i$ and $\sigma_i^a$ ($a = x,y,z$) denotes the three Pauli matrices at sites $i = 1,2,\ldots,N$ of a $d=1$ chain, and $h$ is the transverse field in units of the Ising coupling ($J = 1$ is set in the rest of the paper). We shall mainly consider periodic boundary conditions ($s_N = s_1$), which are the simplest ones allowing for a nonzero energy current to flow through the chain (free boundary conditions imply a zero steady-state current in the above model and as it turns out the same holds for antiperiodic boundary conditions).

The order parameter of this system is $M_x = \sum_i s_i^x$, and the ground state of this Hamiltonian changes from being disordered $\langle M_x \rangle = 0$ for $h > 1$ to ordered $\langle M_x \rangle \neq 0$ for $h < 1$. The transition point $h = h_c$ is a critical point in the universality class of the two-dimensional Ising model.

We would like to investigate the nonequilibrium states of $\hat{H}$ which carry a given energy flux. At zero temperature, this amounts to finding the lowest energy state of $\hat{H}$ with a prescribed energy flux. This can be accomplished [14–17] by introducing a Lagrange multiplier $\lambda$ conjugate to the energy current $\hat{J}_E$. Thus one should find the ground state of the following Hamiltonian:

$$\hat{H} = -J \sum_i s_i^x s_{i+1}^x - \frac{h}{2} \sum_i s_i^z - \lambda \hat{J}_E.$$

(2)

Here the driving field $\lambda$ is again measured in units of $J$, and the current operator $\hat{J}_E$ is the sum of local energy fluxes given by the following expression:

$$\hat{J}_E = \frac{h}{4} \sum_i (s_i^x s_{i+1}^x - s_i^z s_{i+1}^z).$$

(3)

Note that the time evolution of the chain is still governed by $\hat{H}$ through which $\hat{J}_E$ was initially defined. $\hat{H}$ is just a mathematical intermediary to determine the nonequilibrium steady state of $\hat{H}$.

The driven system defined by Eqs. (2) and (3) can be solved [14] and one finds that the ground state does not change and the energy flux is zero up to a critical value $\lambda = \lambda_c(h)$ of the driving field. The ground-state expectation value of energy flux $j_E = \langle \hat{J}_E \rangle / N$ becomes nonzero for $|\lambda| > \lambda_c$. The resulting phase diagram is depicted in Fig. 1.

Here we note a hitherto unnoticed property of the Hamiltonian (2); namely, the duality properties of the transverse Ising model [18] have an appropriate generalization to the full nonequilibrium phase diagram of $\hat{H}$. Indeed, let us denote the action of duality transformation by $s_i^a \rightarrow s_i^{a*}$ and define the transformation as is done for the Ising model in a transverse field

$$s_i^{z*} = 2s_i^x s_{i+1}^x, \quad s_i^{x*} s_{i+1}^{x*} = \frac{1}{2} s_i^z.$$

(4)

It can be verified then that the Hamiltonian $\hat{H}[h,\lambda,\{s_i^a\}]$ (2) characterized by the two couplings $(h,\lambda)$ transforms into an identical Hamiltonian with couplings $(h,\lambda)^* = (h^{-1}, -\lambda h)$,

$$\hat{H}[h,\lambda,\{s_i^a\}] = h \hat{H}[h^{-1}, -\lambda h,\{s_i^{a*}\}].$$

(5)

Hence the duality transformation leaves the whole $|\lambda| = \lambda_c$ curve globally invariant and, furthermore, it leaves the $h = 1$ line pointwise invariant (examples of dual-conjugate points are given in Fig. 1). In order to keep the formulas simple, from now on we shall restrict our analysis to $\lambda \neq 0$, that is to $j_E \neq 0$.

The self-dual $h=1$ line is expected to display special properties. For example, quantities such as $\lambda_c / \lambda$, or the wave vectors $q_\lambda$ where the excitation spectrum is gapless, are left invariant by the duality transformation. Furthermore, the functional form of various physical quantities (dispersion, energy flux, fluctuations) considerably simplify on this line and thus the knowledge of self-duality helps in locating limits where exact calculation can be carried out.

Our main goal is to calculate distribution functions $P(M_x)$ and $P(M_z)$ in various regions in the above phase diagram. These functions are defined as
where brackets denote the quantum mechanical expectation value in the ground state [note that we have omitted the index $\alpha$ in $P_{\alpha}(\cdot)$, i.e., the argument of the function defines which distribution is considered].

There are two parts to our calculations. Function $P(M_{c})$ is evaluated numerically by diagonalizing $\hat{H}$ for chains containing up to $N=18$--20 spins, while $P(M_{c})$ is found analytically. The exact calculation is possible because $M_{c}=\sum_{q}(c_{q}^{\dagger}c_{q}-1/2)$ is a quadratic form in the fermion operators $(c_{q}^{\dagger}, c_{q})$ in which the Hamiltonian $\hat{H}=\hat{H}_{I}-\lambda J_{E}$ is quadratic as well. Thus the calculation of the generating function $G(s)=\langle e^{-sM_{c}}\rangle$ (7)

becomes a problem of evaluating Gaussian integrals. We begin with this part of the problem.

**A. Formalism**

Hamiltonian $\hat{H}$ can be diagonalized [14] by first introducing creation-annihilation operators, then employing the Jordan-Wigner transformation [19] to transform them into fermion operators $(c_{q}, c_{q}^{\dagger})$ and, finally, using a Bogoliubov transformation on the $c_{q}$ and $c_{q}^{\dagger}$ components of the Fourier transforms of $c_{q}$s. The calculation of $P(M_{c})$ becomes relatively simple if after using the Jordan-Wigner transformation one passes to a path-integral formulation (see, e.g., Ref. [20] for a pedagogical account). One then finds that the system is described by the following quadratic action:

$$S[\bar{c}, c] = \int \frac{d\omega}{2\pi} \int \frac{dq}{2\pi} \left[ \frac{1}{2} \bar{c}_{q}(\omega)c_{-q}(-\omega)B_{q,\omega}\left(\bar{c}_{q}(\omega)c_{q}(\omega)\right) \right],$$

where the Grassmann fields $\bar{c}_{q}(\omega)$ and $c_{q}(\omega)$ are related to $c_{q}^{\dagger}$ and $c_{q}$ correspondingly, while the scattering matrix $B_{q,\omega}$ has the following inverse

$$B_{q,\omega}^{-1} = \begin{pmatrix} \langle \bar{c}_{-q}(-\omega)\bar{c}_{q}(\omega) \rangle & \langle \bar{c}_{-q}(-\omega)c_{q}(-\omega) \rangle \\ \langle c_{q}(\omega)\bar{c}_{q}(\omega) \rangle & \langle c_{q}(\omega)c_{q}(-\omega) \rangle \end{pmatrix}.$$ (9)

Here the correlators are given by

$$\langle c_{q}(\omega)\bar{c}_{q}(\omega) \rangle = \langle c_{q}(\omega)c_{q}(\omega) \rangle^*$$

$$= \frac{1}{2\Lambda_{q}} \left[ \frac{h}{2} + \frac{1}{2}\cos q - \Lambda_{q} \right] \left( \frac{1}{2\Lambda_{q}^+} - \frac{1}{i\omega - \Lambda_{q}^+} \right)$$ (10)

and

$$\langle \bar{c}_{-q}(-\omega)\bar{c}_{q}(\omega) \rangle = \langle c_{q}(\omega)c_{-q}(-\omega) \rangle^*$$

$$= \frac{i}{2\Lambda_{q}} \left( \frac{1}{i\omega - \Lambda_{q}^-} - \frac{1}{i\omega - \Lambda_{q}^+} \right).$$ (11)

where $\Lambda_{q}$ and $\Lambda_{q}^\pm$ are the dispersion relations for $\hat{H}_{I}$ and $\hat{H}$, respectively,

$$\Lambda_{q} = \frac{1}{2}\sqrt{1+h^2+2h\cos q},$$ (12)

$$\Lambda_{q}^\pm = \pm \Lambda_{q} + \frac{\lambda h}{4}\sin q.$$ (13)

In order to return to the time variables, we must first study the two branches $\Lambda_{q}^\pm$ of the spectrum.

**B. Spectrum and energy flux**

Flux of energy is present in the system only above a critical drive $\lambda > \lambda_{c}$ [14] (see Fig. 1), where

$$\lambda_{c}(h) = \begin{cases} 2 & \text{if } h \geq 1, \\ \frac{2}{h} & \text{if } h \leq 1. \end{cases}$$ (14)

Indeed, it is not hard to see that

$$\Lambda_{q}^+(\lambda \leq \lambda_{c}) > 0, \quad \Lambda_{q}^-(\lambda \leq \lambda_{c}) < 0.$$ (15)

Hence the ground state is unchanged with respect to the pure Ising model ground state as long as the driving field does not exceed $\lambda_{c}(h)$. This means that all observables will assume their Ising model values and no energy current will be flowing through the chain.

However, for $\lambda \geq \lambda_{c}$ one may see that $\Lambda_{q}^+$ ($\Lambda_{q}^-$) changes sign over the interval $I_{2}\equiv[q_{+}, q_{+}]$ ($I_{4}\equiv[-q_{-}, -q_{-}]$). The explicit expression for the wave vectors $q_{\pm}$ is deduced from

$$\cos q_{\pm} = \frac{-4 \pm \sqrt{(\lambda^2 h^2 - 4)(\lambda^2 - 4)}}{\lambda^2 h}, \quad -\pi \leq q_{\pm} \leq 0.$$ (16)

Beyond the critical drive, the excitation spectrum gives rise to a ground state that breaks the left-right symmetry and, indeed, it may be verified [14] that a nonzero energy flux is present in the chain with the explicit form of the flux given by

$$f_{E} = \frac{h}{4\pi \lambda^2} \sqrt{(\lambda^2 - \frac{4}{h^2})(\lambda^2 - 4)}.$$ (17)
After normal ordering, we further define here intervals

\[ I = \begin{cases} \lambda - 2, & \text{for } h > 1, \ \lambda > 2, \\ 1/4 \pi, & \text{for } h = 1, \ \lambda = 2, \\ \sqrt{1 - h^2} / 4 \pi, & \text{for } h < 1, \ \lambda = 2. \]

Besides, as they will naturally arise in the upcoming discussion, we further define here intervals \( I_1 = [-\pi, q_+] \), \( I_2 = [-q_+, \pi] \), and \( I_3 = [-q_-, \pi] \) which are complementary to \( I_2 \) and \( I_4 \) in \([-\pi, \pi]\).

C. Inverse of the scattering matrix in the absence or presence of energy flux

When no current flows through the system, \( j_x = 0 \), the equal-time transform of \( (B_{q, \omega})^{-1} \), denoted by \( B_{q}^{-1} \), reads

\[
B_{q}^{-1} = \begin{pmatrix}
\frac{1}{2} \sin q & \frac{2 h}{2 + \cos q + \Lambda_q} \\
\frac{4 \Lambda_q}{2 + \cos q + \Lambda_q} & \frac{2 h}{2 + \cos q + \Lambda_q} \\
\frac{2 h}{2 + \cos q + \Lambda_q} & \frac{1}{2} \sin q
\end{pmatrix}.
\]

For \( j_x \neq 0 \), on the other hand, \( B_{q}^{-1} \) is given by Eq. (19) only for \( q \in I_1 \cup I_2 \cup I_3 \). If \( q \in I_2 \cup I_4 \), its expression is changed to

\[
I_2: B_{q}^{-1} = \begin{pmatrix}
0 & 0 \\
0 & 0
\end{pmatrix}, \quad I_4: B_{q}^{-1} = \begin{pmatrix}
0 & 1 \\
0 & 0
\end{pmatrix}.
\]

Having the expressions for \( B_{q}^{-1} \), we can start the calculation of the distribution function \( P(M_z) \).

III. DISTRIBUTION FUNCTION FOR THE TRANSVERSE MAGNETIZATION

A. Calculation of the generating function and its moments

The generating function \( G(s) \) of \( P(M_z) \) can be expressed through fermionic operators as

\[
G(s) = \langle e^{-sM_z} \rangle = e^{Ns/2} \left\langle \exp \left(-s \sum q \bar{c}_q c_q\right) \right\rangle.
\]

After normal ordering \( e^{-s \bar{c}_q c_q} \) and using the Grassmann fields, we are left with evaluating the following expression:

\[
G(s) = e^{Ns/2} \left\langle \exp \left(-s \sum q \bar{c}_q c_q\right) \right\rangle, \quad s = 1 - e^{-s}
\]

and the fields in the exponentials are evaluated at some fixed time. In order to evaluate Eq. (22), we recall the following result for Grassmann integrals [20]: the vacuum expectation value of observables of form \( \exp(-\sum q \bar{c}_q c_q) \) can be obtained as

\[
\langle \exp(-s \sum q \bar{c}_q c_q) \rangle = \prod_q \sqrt{\det \left[ -\frac{1}{2} \Lambda_q \right]}.
\]

As one can see \( G(s) \) is a special case of Eq. (23) and it can be evaluated by using the appropriate expressions (19) or (20) for \( B_{q}^{-1} \).

If no current flows, that is, for \( \lambda \ll \lambda_c \), we find that \( \ln G \) (the cumulant generating function) is given by

\[
\ln G(s) = \frac{N}{2} \int_{-\pi/2}^{\pi/2} dq (1 - n_q) (e^{s} + n_q e^{-s}),
\]

where

\[
n_q = (h + \cos q + 2 \Lambda_q) / (4 \Lambda_q).
\]

One can verify that the normalization condition \( G(0) = 1 \) is satisfied, and one can also recover the well-known result found by Pfeuty [21] for the magnetization

\[
-\frac{\partial \ln G(s)}{\partial s} \bigg|_{s=0} = \langle M_z \rangle = N \int_{-\pi/2}^{\pi/2} dq (n_q - 1/2).
\]

As we shall see below, \( P(M_z) \) is a Gaussian thus, in addition to \( \langle M_z \rangle \), the variance of the transverse magnetization \( N \sigma^2(h) = \langle M_z^2 \rangle - \langle M_z \rangle^2 \) will characterize the distribution. It can be obtained from the second derivative of \( \ln G(s) \) as

\[
\sigma^2(h) = 2 \int_{-\pi/2}^{\pi/2} dq n_q (1 - n_q).
\]

It is interesting to note that the fluctuations in \( M_z \) are independent of the magnetic field in the ordered phase

\[
\sigma^2(h) = \begin{cases} 1/4 & \text{for } |h| \leq 1, \\
1/4 h^2 & \text{for } |h| > 1.
\end{cases}
\]

In the presence of nonzero energy flux \( \lambda \gg \lambda_c \), the generating function is more complicated only because of the limits of integration in Eq. (24).

\[
\ln G_{\lambda}(s) = \frac{N}{2} \int_{I_1 \cup I_3 \cup I_4} dq \ln (1 - n_q) \left( e^{s} + n_q e^{-s} \right).
\]

Accordingly, the first and second cumulants of the magnetization are given by

\[
\langle M_z \rangle_{\lambda} = \frac{N}{2} \int_{I_1 \cup I_3 \cup I_4} dq \frac{h + \cos q}{2 \Lambda_q},
\]

\[056129-4\]
\[ \sigma^2_s(h) = \frac{1}{2} \int_{t_1 \cup t_2 \cup t_3} dq \frac{\sin^2 q}{1 + h^2 + 2h \cos q}, \]  

(31)

where we have used Eq. (25) to write out the integrands explicitly. Note that we added a \( \lambda \) subscript to \( G \), \( \langle M^2 \rangle_c \), and \( \sigma^2 \) in order to indicate that these quantities do depend on \( \lambda \) for \( \lambda > \lambda_c \).

**B. The transverse magnetization distribution is a Gaussian**

In order to show that the transverse magnetization distribution is a Gaussian, let us consider the \( n \)th cumulant \( \langle M^n \rangle_c \) which is the coefficient in front of \( (-s)^n/n! \) in the expansion of \( \ln G(s) \). The latter coefficient is \( N \) times an integral of a polynomial in \( n_q \), both in the current-carrying and current-free phases and, furthermore, \( n_q \) is a nonsingular, strictly positive function of \( q \) (note that \( n_x \) is finite even at \( h = 1 \)). Hence each cumulant depends linearly on \( N \). In particular, as we have seen, the variance of \( M_c \) denoted by \( N\sigma^2 \) and \( \sigma^2 \) [in Eq. (31) or Eq. (27)] is finite.

It follows from the linear \( N \) dependence of the cumulants that

\[ \frac{\langle M^n \rangle_c}{\langle M_c \rangle_c^{n/2}} \sim \frac{N}{N^{n/2}} \sim N(2-n)/2, \]

(32)

thus the above ratio goes to zero for all \( n > 2 \) as \( N \to \infty \). We may therefore conclude that the limiting form of the distribution function of the transverse magnetization is a Gaussian of variance \( N\sigma^2 \):

\[ P(M_c) = \frac{1}{\sqrt{2\pi N\sigma^2}} e^{-(M_c - \langle M_c \rangle)^2/2N\sigma^2}. \]

(33)

The above result applies over the whole phase diagram and it is useful to check the numerical procedure employed in Sec. IV by evaluating \( P(M_c) \) for finite-size systems. As can be seen in Fig. 2, there is a convergence to the limiting form with increasing \( N \) and, furthermore, the nearly Gaussian fluctuations of \( M_c \) are observed already at small sizes \( (N = 16 - 20) \). It is remarkable that the deviations from the Gaussian are very close to those of an \( N \) step random walk with a drift determined from a correspondence between the left (right) moves and the up (down) spins generating the average \( \langle M_c \rangle \).

**C. Width of the Gaussian**

Recalling the expression of the average and the variance, Eqs. (26) and (27), we calculate them in the limit of vanishing flux \( (\lambda \to \lambda_c^+) \). Let us define

\[ \delta M_c = \langle M_c \rangle_c - \langle M_c \rangle, \quad \delta \sigma^2 = \sigma^2(h) - \sigma^2(h). \]

(34)

The above quantities exhibit singular behavior as one enters the current-carrying phase. For the magnetization we find

\[ \delta M_c = -N \int_{q_+}^{q_-} dq \frac{h + \cos q}{2\pi} \Lambda_q, \]

(35)

**IV. DISTRIBUTION FUNCTION FOR THE LONGITUDINAL MAGNETIZATION**

The exact evaluation of \( P(M_x) \) appears to be a nontrivial task and we have been able to calculate it only numerically for finite-size chains. Since expression (6) for \( P(M_x) \) is a ground-state expectation value, we had to find the ground-state wave function and, due to the sparseness of the Hamiltonian matrix, the Lanczos algorithm [22] could be used ef-
1, and a nontrivial distribution emerges only at criticality (to be a Gaussian for chain lengths of up to ∞nite. Thus one expects effectively. Since the ground-state wave function is needed Figs. 3–5.

A. Equilibrium distribution

In the equilibrium system (J_E=0), the correlation length is infinite only at the critical point. Thus one expects P(M_x) to be a Gaussian for h>1, a sum of two Gaussians for h<1, and a nontrivial distribution emerges only at criticality (h=h_c=1). This is indeed what we observe, apart from the finite-size effects showing up in non-Gaussian corrections close to h=1. At the critical point itself, P(M_x) shows fast convergence to the asymptotic form as can be seen in Fig. 3 where the N=16 points appear to have settled on the asymptotic curve. This means that the N dependence of P(M_x) is almost all in the scaling variable M_x/√⟨M_x^2⟩.

B. Nonequilibrium distribution

In the nonequilibrium case (J_E ≠ 0), we find that similarly to M_c, the fluctuations of the longitudinal magnetization M_x are also Gaussian (Fig. 5). Remarkably, the finite-size, non-Gaussian corrections are small even near the h=1, λ=2 point.

The Gaussian result is somewhat surprising since the flux generates power-law correlations in ⟨s_{n+1}^a s_n^a⟩ decaying with distance as 1/√n [14]. Thus one may presume that the current-carrying states are effectively critical. This is not so, however, since the correlations are oscillating and the fluctuations ⟨M_x^2⟩ (given by the integral of the correlations) are finite. Actually ⟨M_x^2⟩ is decreasing with increasing J_E (see...
The decrease of $\langle M_1^2 \rangle$ with increasing $j_E$ can be seen in the numerical studies of the finite spin chains. On the self-dual line ($h = 1$), however, this can be shown analytically in the limit of vanishing flux ($\lambda \rightarrow \lambda_c = 2$), where one finds

$$\langle M_1^2 \rangle \sim Nj_E^{-3/8}. \quad (39)$$

The derivation of the above result is possible because the ($h = 1$, $\lambda = \lambda_c$) point is a critical point with infinite correlation length. Approaching this point along the $h = 1$, $\lambda \rightarrow \lambda_c^+$ line, one can observe from numerical studies [14] that the wavelength $\kappa^{-1}$ of the oscillation of the correlation function $C_s(r) = \langle s_i s_{i+r} \rangle \sim \cos kr$ diverges as $\kappa^{-1} \sim (\lambda - \lambda_c)^{-1/2}$. This diverging wavelength allows one to take a continuum limit and to establish the order-parameter correlations possess the following scaling form:

$$C_s(r) = \frac{A}{r^{1/4}} F(kr), \quad (40)$$

with the small argument limit of the scaling function explicitly given by $F(x) = 1 - x^2/2 + x^4/16 + O(x^5)$.

The derivation of the above results follows the ideas used for the calculation of the order-parameter fluctuations in the equilibrium transverse Ising model; namely, the correlations are expressed as a Pfaffian of a block Toeplitz matrix constructed of $2 \times 2$ matrices. Then, in the equilibrium case, the analysis of the Toeplitz matrices in the asymptotic limit of $r \rightarrow \infty$, and $h \rightarrow 1$ with $r(h-1)$ kept fixed, yields $C_s(r) \sim r^{-1/4} F(r(h-1))$. A similar asymptotic analysis in the current-carrying phase, using the scaling limit $r \rightarrow \infty$, and $\kappa^{-1} \sim (\lambda - \lambda_c)^{1/2} \rightarrow 0$ with $kr$ kept finite, results in Eq. (40). The derivation is rather technical and we shall present it, together with the analysis of other scaling limits, in a separate publication [33].

Once the correlations are known, the fluctuations can be calculated from

$$\langle M_1^2 \rangle = N \left[ 1 + 2 \sum_{r \geq 1} C_s(r) \right] \approx N \int_0^{+\infty} dr C_s(r), \quad (41)$$

where changing the sum into integral is again allowed because of the diverging characteristic length scale $\kappa^{-1}$. Using now the scaling form (40), we find that

$$\langle M_1^2 \rangle \sim N \kappa^{-3/4} \sim N j_E^{-3/8}. \quad (42)$$

The above expression demonstrates the decrease of fluctuations with increasing flux and it also tells us how the Gaussian distribution crosses over to the nontrivial shape observed at the critical point.

V. FINAL REMARKS

Returning to the problems discussed in the Introduction, we can see that the connection between NESS and critical states in terms of (universal) distribution functions is not straightforward. NESS is generated by fluxes, and fluxes may or may not generate long-range correlations. There are numerous examples [1–3] where the fluxes are spatially localized and long-range correlations do not develop (unless the system is at a special point in the parameter space). Clearly, in such cases, one cannot hope for a general description to emerge. If the fluxes are global, as is the case for the model treated in the present paper, long-range correlations do emerge frequently [1–3]. Even in this case, however, it is far from trivial whether these correlations drive the system to an effectively critical state or whether they make the system more rigid.

The driven transverse Ising model treated above is an example where a global flux of energy generates long-range correlations but the resulting state becomes more rigid in the sense that the fluctuations are decreased due to the presence of the flux. Driven diffusive systems [1] provide other examples [34] where the fluctuations decrease while the fluxes induce power-law correlations. Thus we should conclude that, in general, the power-law correlations generated by global fluxes cannot be the source of possible universality of nonequilibrium distribution functions. It remains, however, an intriguing question whether the weak long-range interactions supported by global fluxes can underlie a kind of “weak” universality classification of distributions in NESS.

ACKNOWLEDGMENTS

We thank G. Györgyi, V. Hunyadi, and L. Sasvári for helpful discussions. This research was partially supported by the Hungarian Academy of Sciences (Grant No. OTKA T029792).


