

Spectrum of the density matrix of a large block of spins of the XY model in one dimension

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Abstract We consider reduced density matrix of a large block of consecutive spins in the ground states of the XY spin chain on an infinite lattice. We derive the spectrum of the density matrix using expression of Rényi entropy in terms of modular functions. The eigenvalues λ_n form exact geometric sequence. For example, for strong magnetic field $\lambda_n = C \exp(-\pi \tau_0 n)$, here $\tau_0 > 0$ and $C > 0$ depend on anisotropy and magnetic field. Different eigenvalues are degenerated differently. The largest eigenvalue

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is unique, but degeneracy g_n increases sub-exponentially as eigenvalues diminish: $g_n \sim \exp(\pi\sqrt{n/3})$. For weak magnetic field expressions are similar.

Keywords Entanglement · Entropy · Entanglement spectrum · Reduced density matrix · XY model · Partitions theory

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1 Introduction

Entanglement distinguishes quantum systems from classical ones. While the notion of entanglement has been introduced since the dawn of quantum mechanics, only recently physicists have employed a quantitative approach, mostly inspired by the progresses in quantum information and Bethe ansatz.

The most studied quantity is the *Von Neumann entropy*, which is the quantum analog of the Shannon entropy. It is a reasonable measure of entanglement for pure states. The entropy in a system A described by a density matrix ρ_A :

$$S(\rho_A) \equiv -\text{Tr}(\rho_A \ln \rho_A). \quad (1.1)$$

We shall consider simplest case of a pure system U described by a wave function $|\Psi\rangle$. The system U is composed by the union of two disjoint subsystems A and B , meaning $U \equiv A \cup B$. The reduced density matrix of subsystem A obtained by tracing out the B degrees of freedom $\rho_A \equiv \text{Tr}_B |\Psi\rangle\langle\Psi|$ and the Von Neumann entropy of a subsystem is a measure of entanglement between the two subsystems¹ [21].

Importance of *Renyi entropy*

$$S_R(\rho_A, \alpha) \equiv \frac{1}{1-\alpha} \ln \text{Tr}(\rho_A^\alpha), \quad 1 > \alpha > 0. \quad (1.2)$$

for quantum information was emphasized in [1, 8, 10, 25] Note that in the limit $\alpha \rightarrow 1$, the Renyi entropy recovers the Von Neumann entropy:

$$\lim_{\alpha \rightarrow 1} S_R(\rho_A, \alpha) = S(\rho_A).$$

We calculated the *Renyi entropy* of large block of spins for XY spin chain [14, 18] and studied its analytical continuation into the complex plane of parameter α .

Another quantity that has recently attracted a lot of interest is the spectrum of the reduced density matrix. The knowledge of the entanglement spectrum fixes the density matrix up to unitarity transformation, we can say that it fixes the state of the block.

¹ Note that $S(\rho_A) = S(\rho_B)$.

It is also clear that the Renyi entropy and the entanglement spectrum are also very closely related. In fact, if we know ζ -function of ρ_A is:

$$\zeta_{\rho_A}(\alpha) \equiv \text{tr} \rho_A^\alpha = \sum_{n=0}^{\infty} g_n \lambda_n^\alpha. \tag{1.3}$$

Originally the Renyi entropy is defined for $0 < \alpha < 1$, but analytical expression which we derived for it helped to construct analytic continuation to the whole complex plane of α . At all α then we can find eigenvalues λ_n , which have the meaning of probabilities ($0 < \lambda_n < 1$) and their multiplicities g_n . From the definition of the Renyi entropy (1.2) we see that

$$\zeta_{\rho_A}(\alpha) = e^{(1-\alpha)S_R(\rho_A, \alpha)}. \tag{1.4}$$

The aim of this paper is to use this relation to calculate the spectrum of ρ_A for the anisotropic XY model, for which analytical expressions for the Renyi entropy are known explicitly [14, 18]. The XY model is one of the simplest integrable models (since quasiparticle excitations are essentially free fermions), while still having an interesting and non-trivial phase diagram. Our approach is based on the paper [27], where Von Neumann entropy of block of spins in the ground state of the XY model was represented as a determinant of Toeplitz matrix. We used Riemann-Hilbert problem to evaluate the Toeplitz determinant for large block of spins [15]. Analytical properties of Von Neumann entropy was studied in [13]. The method of [15] was generalized in [16, 17, 19] for more general quantum spins. The evaluation of Renyi Entropy for the XY model, which was done in [14, 18], uses again the Fisher-Hartwig formulae and the Riemann-Hilbert approach. Although this is the first time that these exact results are used to access the full spectrum of ρ_A , general behaviors has already been known, because of the underlying free fermionic structure.

The degeneracy of the eigenvalues is given by the number of ways a given energy can be realized by different excitations. This is essentially a partitioning problem and it has been addressed already several years ago in an effort toward the optimization of Density Matrix Renormalization Group approaches [22]. Our exact approach agrees with results of [22].

It is also should be mentioned, that for some spin models reduced density matrix was evaluated explicitly [26, 28].

2 Quantum entropies for the XY model

The anisotropic XY spin chain is defined by the Hamiltonian

$$H = - \sum_{j=-\infty}^{\infty} \left[(1 + \gamma) \sigma_j^x \sigma_{j+1}^x + (1 - \gamma) \sigma_j^y \sigma_{j+1}^y + h \sigma_j^z \right], \tag{2.1}$$

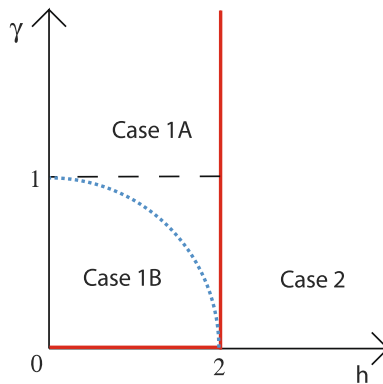


Fig. 1 Phase diagram of the anisotropic XY model in a constant magnetic field (only $\gamma \geq 0$ and $h \geq 0$ shown). The three cases 2, 1A, 1B, considered in this paper, are clearly marked. The critical phases ($\gamma = 0, h \leq 2$ and $h = 2$) are drawn in *bold lines (red)*. The boundary between cases 1A and 1B, where the ground state is given by two degenerate product states, is shown as a *dotted line (blue)*. The Ising case ($\gamma = 1$) is also indicated, as a *dashed line*. (color figure online)

where γ is the anisotropy parameter, σ_j^x, σ_j^y and σ_j^z are the Pauli matrices and h is the magnetic field. The Hamiltonian is clearly symmetric under the transformations $\gamma \rightarrow -\gamma$ or $h \rightarrow -h$, therefore we can consider just the quadrant $\gamma > 0$ and $h \geq 0$ without loss of generality. The system is gapped in the bulk of the phase diagram and has two phase transitions where the spectrum becomes critical: at $\gamma = 0, |h| < 2$ one has the XX model (universality of free fermions on a lattice) and $|h| = 2$ is the critical magnetic field of the Ising phase transition, see Fig. 1.

The model was solved in [2–4, 6, 7, 20]. It is known that its correlation functions can be calculated using methods of Toeplitz determinants and Riemann-Hilbert problems. These techniques were applied in [13–15, 18] to the study of the quantum entropies.

The density matrix of the unique ground state $|GS\rangle$ of the model is given by $\rho_{AB} = |GS\rangle\langle GS|$. The reduced density matrix of a subsystem A is $\rho_A = Tr_B(\rho_{AB})$. We take the subsystem A to be a block of d consecutive spins (system B is the state of the rest of the chain) and consider the double scaling limit $1 \ll d \ll N = \infty$, where N is the total number of sites in the chain, which we take to be infinite. The size of the block is very large, but still is much smaller than the total size of the lattice. Matrix elements of ρ_A are correlation functions, see formulae (17) and (18) of [18].

The Rényi entropy (1.2) converges to the von Neumann entropy (1.1) for $\alpha \rightarrow 1$, therefore we can concentrate just on the former quantity. Its analytical expressions were derived in [14] and they can be written as

$$S_R(\rho_A, \alpha) = \frac{1}{6} \frac{\alpha}{1-\alpha} \ln(k' k) + \frac{1}{3} \frac{1}{1-\alpha} \ln \left(\frac{\theta_3^2(0|\alpha i \tau_0)}{\theta_2(0|\alpha i \tau_0)\theta_4(0|\alpha i \tau_0)} \right) + \frac{1}{3} \ln 2, \tag{2.2}$$

for $h > 2$ and

$$S_R(\rho_A, \alpha) = \frac{1}{6} \frac{\alpha}{1-\alpha} \ln\left(\frac{k'}{k^2}\right) + \frac{1}{3} \frac{1}{1-\alpha} \ln\left(\frac{\theta_2^2(0|\alpha i \tau_0)}{\theta_3(0|\alpha i \tau_0)\theta_4(0|\alpha i \tau_0)}\right) + \frac{1}{3} \ln 2, \tag{2.3}$$

for $h < 2$; where

$$\tau_0 \equiv \frac{I(k')}{I(k)}, \quad k' = \sqrt{1-k^2}, \tag{2.4}$$

$I(k)$ is the complete elliptic integral of the first kind,

$$I(k) = \int_0^1 \frac{dx}{\sqrt{(1-x^2)(1-k^2x^2)}} \tag{2.5}$$

and

$$\theta_j(z|\tau) := \theta_j(z, q), \quad q = e^{\pi i \tau} \quad j = 1, 2, 3, 4, \tag{2.6}$$

are the elliptic theta functions defined by the following Fourier series ($|q| < 1$)

$$\theta_1(z, q) = i \sum_{m=-\infty}^{\infty} (-1)^m q^{\binom{2m-1}{2}} e^{2iz(m-\frac{1}{2})}, \tag{2.7}$$

$$\theta_2(z, q) = \sum_{m=-\infty}^{\infty} q^{\binom{2m-1}{2}} e^{2iz(m-\frac{1}{2})}, \tag{2.8}$$

$$\theta_3(z, q) = \sum_{m=-\infty}^{\infty} q^{m^2} e^{2izm}, \tag{2.9}$$

$$\theta_4(z, q) = \sum_{m=-\infty}^{\infty} (-1)^m q^{m^2} e^{2izm}. \tag{2.10}$$

The elliptic parameter $k = k(\gamma, h)$ is defined in the different regions of the phase diagram as

$$k \equiv \begin{cases} \sqrt{(h/2)^2 + \gamma^2 - 1} / \gamma, & \text{Case 1a : } 4(1 - \gamma^2) < h^2 < 4; \\ \sqrt{(1 - h^2/4 - \gamma^2)/(1 - h^2/4)}, & \text{Case 1b : } h^2 < 4(1 - \gamma^2); \\ \gamma / \sqrt{(h/2)^2 + \gamma^2 - 1}, & \text{Case 2 : } h > 2. \end{cases} \tag{2.11}$$

Alternatively, we can write the Rényi in terms of the λ -modular function (see [14]) as

$$S_R = \begin{cases} \frac{1}{6} \frac{\alpha}{1-\alpha} \ln(k k') - \frac{1}{12} \frac{1}{1-\alpha} \ln[\lambda(\alpha i\tau_0) (1 - \lambda(\alpha i\tau_0))] + \frac{1}{3} \ln 2 & h > 2 \\ \frac{1}{6} \frac{\alpha}{1-\alpha} \ln\left(\frac{k'}{k^2}\right) - \frac{1}{12} \frac{1}{1-\alpha} \ln\left[\frac{1-\lambda(\alpha i\tau_0)}{\lambda^2(\alpha i\tau_0)}\right] + \frac{1}{3} \ln 2 & h < 2 \end{cases} \quad (2.12)$$

The modular function is defined as ($\Im\tau > 0$)

$$\lambda(\tau) = \frac{\theta_2^4(0|\tau)}{\theta_3^4(0|\tau)} \equiv k^2(\tau), \text{ and } 1 - \lambda(\tau) = \frac{\theta_4^4(0|\tau)}{\theta_3^4(0|\tau)} \equiv k'^2(\tau), \quad (2.13)$$

We note that the basic modular parameter $k \equiv k(\gamma, h)$ defined in (2.11) coincide with the value of the function $k(\tau)$ at $\tau = i\tau_0$.

A third representation of the Renyi entropy in terms of q -series will be useful to determine the multiplicities of the reduced density matrix eigenvalues [14]:

$$S_R(\rho_A, \alpha) = \begin{cases} \frac{1}{12} \frac{\alpha}{1-\alpha} \ln\left(\frac{k^2 k'^2}{16q}\right) + \frac{2}{1-\alpha} \ln \prod_{m=0}^{\infty} [1 + q^{(2m+1)\alpha}] & h > 2 \\ \frac{1}{6} \frac{\alpha}{1-\alpha} \ln\left(\frac{16q k'}{k^2}\right) + \frac{2}{1-\alpha} \ln \prod_{m=1}^{\infty} [1 + q^{2m\alpha}] + \ln 2 & h < 2 \end{cases} \quad (2.14)$$

where

$$q \equiv e^{-\pi\tau_0} = e^{-\pi I(k')/I(k)}. \quad (2.15)$$

3 Spectrum of ρ_A

Using the expressions for the Renyi entropy we just listed, we now want to determine the eigenvalues λ_n ($0 < \lambda_n < 1$) the operator ρ_A and their multiplicities g_n , through its momentum function (1.3) using (1.4). Using (2.14) we have

$$\zeta_{\rho_A}(\alpha) = \begin{cases} e^{\alpha\left(\frac{\pi\tau_0}{12} + \frac{1}{6} \ln \frac{k k'}{4}\right)} \prod_{m=0}^{\infty} (1 + q_\alpha^{2m+1})^2 & h > 2 \\ 2e^{\alpha\left(-\frac{\pi\tau_0}{6} + \frac{1}{6} \ln \frac{k'}{4k^2}\right)} \prod_{m=1}^{\infty} (1 + q_\alpha^{2m})^2 & h < 2 \end{cases}, \quad (3.1)$$

where

$$q_\alpha \equiv e^{-\alpha\pi\tau_0} = q^\alpha. \quad (3.2)$$

To use these expression, we will need some results on q -series and elementary notions of the theory of partitions.

Let us concentrate first on the case $h > 2$. Classical arguments of the theory of partitions (see e.g. [5]) tell us that

$$\prod_{n=0}^{\infty} (1 + q^{2n+1}) = \sum_{n=0}^{\infty} p_{\mathcal{O}}^{(1)}(n)q^n, \tag{3.3}$$

where $p_{\mathcal{O}}^{(1)}(0) = 1$ and $p_{\mathcal{O}}^{(1)}(n)$, for $n > 1$, denotes the number of partitions of n into distinct positive *odd* integers, i.e.

$$p_{\mathcal{O}}^{(1)}(n) \equiv \#\{(m_1, \dots, m_k) : m_j = 2r_j + 1, \quad m_1 > m_2 > \dots > m_k, \quad n = m_1 + m_2 + \dots + m_k\}. \tag{3.4}$$

Hence (3.1) for $h > 2$ becomes

$$\zeta_{\rho_A}(\alpha) = e^{\alpha\left(\frac{\pi\tau_0}{12} + \frac{1}{6} \ln \frac{k k'}{4}\right)} \sum_{n=0}^{\infty} a_n q_{\alpha}^n, \tag{3.5}$$

where,

$$a_0 = 1, \quad a_n = \sum_{l=0}^n p_{\mathcal{O}}^{(1)}(l) p_{\mathcal{O}}^{(1)}(n-l) \tag{3.6}$$

Since

$$q_{\alpha}^n = (e^{-\pi\tau_0 n})^{\alpha}, \tag{3.7}$$

we conclude that

$$\zeta_{\rho_A}(\alpha) = \sum_{n=0}^{\infty} a_n \lambda_n^{\alpha}, \quad \lambda_n = e^{-\pi\tau_0 n + \frac{\pi\tau_0}{12} + \frac{1}{6} \ln \frac{k k'}{4}}. \tag{3.8}$$

Comparing the last equation with (1.3) we arrive at the following theorem:

Theorem 3.1 *Let the magnetic field $h > 2$. Then, the eigenvalues of the reduced density matrix ρ_A are given by*

$$\lambda_n = e^{\frac{1}{6} \ln \frac{k k'}{4} - \pi \frac{l(k')}{l(k)} \left[n - \frac{1}{12}\right]}, \quad n = 0, 1, 2, \dots \tag{3.9}$$

and the corresponding multiplicities $g_n = a_n$ are defined by (3.6). Used notations are (2.5), (2.11) and (2.4).

The case $h < 2$ is treated in a very similar way. Instead of (3.3) we use another combinatorial identity,

$$\prod_{n=1}^{\infty} (1 + q^{2n}) = \sum_{n=0}^{\infty} p_{\mathcal{N}}^{(1)}(n)q^{2n}, \tag{3.10}$$

where $p_{\mathcal{N}}^{(1)}(0) = 1$ and $p_{\mathcal{N}}^{(1)}(n)$, for $n > 1$, denotes the number of partitions of n into distinct positive integers, i.e.

$$p_{\mathcal{N}}^{(1)}(n) \equiv \#\{(m_1, \dots, m_k) : m_1 > m_2 > \dots > m_k > 0, n = m_1 + m_2 + \dots + m_k\}. \tag{3.11}$$

It is worth noticing that (see e.g. [5])

$$p_{\mathcal{N}}^{(1)}(n) = p_{\mathcal{O}}(n), \tag{3.12}$$

where $p_{\mathcal{O}}(n)$ denotes the partitions of n into positive *odd* integers:

$$p_{\mathcal{O}}(n) \equiv \#\{(m_1, \dots, m_k) : m_j = 2r_j + 1, n = m_1 + m_2 + \dots + m_k\}. \tag{3.13}$$

Here we assume that m_k form an increasing sequence. The analog of Eq. (3.5), with the help of (3.10), now reads as

$$\begin{aligned} \zeta_{\rho_A}(\alpha) &= 2e^{\alpha\left(-\frac{\pi\tau_0}{6} + \frac{1}{6} \ln \frac{k'}{4k^2}\right)} \sum_{n=0}^{\infty} b_n q_{\alpha}^{2n} \\ &= 2 \sum_{n=0}^{\infty} b_n \lambda_n^{\alpha}, \quad \lambda_n = e^{-2\pi\tau_0 n - \frac{\pi\tau_0}{6} + \frac{1}{6} \ln \frac{k'}{4k^2}}, \end{aligned} \tag{3.14}$$

where

$$b_0 = 1, \quad b_n = \sum_{l=0}^n p_{\mathcal{N}}^{(1)}(l) p_{\mathcal{N}}^{(1)}(n-l). \tag{3.15}$$

Finally, comparing (3.14) with Eq. (1.3) we arrive at the analog of Theorem 3.1 for the case $h < 2$:

Theorem 3.2 *Let the magnetic field $h < 2$. Then, the eigenvalues of the reduced density matrix ρ_A are given by the equation,*

$$\lambda_n = e^{\frac{1}{6} \ln \frac{k'}{4k^2} - 2\pi \frac{l(k')}{l(k)} \left[n + \frac{1}{12}\right]}, \quad n = 0, 1, 2, \dots \tag{3.16}$$

and the corresponding multiplicities $g_n = 2b_n$ where the integers b_n are determined by (3.15). Used notations are (2.5), (2.11) and (2.4).

4 Asymptotics of g_n

Consider first the case $h > 2$. Following the usual methodology, we introduce the generating function

$$f(z) := \sum_{n=0}^{\infty} g_n z^n. \tag{4.1}$$

This function is holomorphic in the unit disc, $\tau_0 > 0$. Indeed, we have from (3.5) that,

$$f(z) = e^{-\alpha \left(\frac{1}{6} \ln \frac{k k'}{4} + \frac{\pi \tau_0}{12} \right)} \zeta_{\rho_A}(\alpha), \quad \alpha = -\frac{1}{\pi \tau_0} \ln z. \tag{4.2}$$

Statement of holomorphicity then follows from the first equation in (3.1).

The function $f(z)$ has a singularity at $z = 1$. In order to see this, we deduce from (1.4) the representation for $f(z)$ in terms of the entropy $S_R(\rho_A, \alpha)$

$$f(z) = e^{(1-\alpha)S_R(\rho_A, \alpha) - \alpha \left(\frac{1}{6} \ln \frac{k k'}{4} + \frac{\pi \tau_0}{12} \right)}, \quad \alpha = -\frac{1}{\pi \tau_0} \ln z. \tag{4.3}$$

In [14], using the explicit formulae (2.2) and (2.12), and the modular properties of the λ -function,

$$\lambda \left(-\frac{1}{\tau} \right) = 1 - \lambda(\tau), \quad \lambda(\tau + 1) = \frac{\lambda(\tau)}{\lambda(\tau) - 1}, \tag{4.4}$$

it was obtained that

$$S_R(\rho_A, \alpha) = \frac{1}{\alpha(1-\alpha)} \frac{\pi}{12} \frac{I(k)}{I(k')} + \frac{\alpha}{1-\alpha} \frac{1}{6} \ln \frac{k k'}{4} + O \left(e^{-\frac{\pi}{\alpha \tau_0}} \right),$$

$$\alpha \rightarrow 0, \quad -\frac{\pi}{2} < \arg \alpha < \frac{\pi}{2}. \tag{4.5}$$

Hence,

$$f(z) = e^{-\frac{\pi^2}{12 \ln z} + \frac{1}{12} \ln z + O(e^{\pi^2 / \ln z})}, \quad z \rightarrow 1, \quad |z| < 1. \tag{4.6}$$

The coefficients g_n are given by the Cauchy formula,

$$g_n = \frac{1}{2\pi i} \int_{|z|=1-\epsilon} \frac{f(z)}{z^{n+1}} dz, \tag{4.7}$$

From this formula, the large n asymptotics of g_n can be rigorously obtained by Hardy-Ramanujan-Rademacher circle method (see e.g. [24]) using modular properties (4.4) of the λ -function. According to the circle method, the leading contribution to integral (4.7)

comes from the neighborhood² of the point $1 - \epsilon$. This fact can be also demonstrated by plotting the function $g(z) = \ln f(z) - n \ln z$, see Fig. 2. Therefore, we can replace the explicit formula (4.7) by the estimate,

$$g_n \simeq \frac{1}{2\pi i} \int_{\mathcal{L}} \frac{f(z)}{z^{n+1}} dz \simeq \frac{1}{2\pi i} \int_{\mathcal{L}} \frac{e^{-G(z)}}{z} dz, \quad n \rightarrow \infty, \tag{4.8}$$

where $\mathcal{L} = \{|z| = \rho_n, |\arg z| < \delta \leq -\ln \rho_n\}$, and

$$G(z) \equiv \frac{\pi^2}{12 \ln z} + \left(n - \frac{1}{12}\right) \ln z. \tag{4.9}$$

We determine ρ_n as the stationary point of $G(z)$, i.e.:

$$\left. \frac{dG(z)}{dz} \right|_{z=\rho_n} = 0 \Rightarrow \rho_n = e^{-\frac{\pi}{\sqrt{12n-1}}} \equiv e^{-\epsilon_n}. \tag{4.10}$$

Switching then to polar coordinate $z = e^{-\epsilon_n + i\theta}$, we can rewrite $G(z)$ in the form,

$$\begin{aligned} G(z) &= \frac{\pi^2}{12} (-\epsilon_n + i\theta)^{-1} + \left(n - \frac{1}{12}\right) (-\epsilon_n + i\theta) \\ &= -\frac{\pi}{\sqrt{3}} \sqrt{n - \frac{1}{12}} + \frac{\pi^2}{12\epsilon_n^3} \theta^2 \left(1 + O\left(\frac{\theta}{\epsilon_n}\right)\right), \quad z \in \mathcal{L}. \end{aligned} \tag{4.11}$$

It can be shown, using again the circle method, that there exists a positive κ_0 such that the following choice of the parameter δ in the definition of the arc \mathcal{L} is consistence with estimate (4.8):

$$\delta = \epsilon_n^{1+2\kappa}, \quad 0 < \kappa < \kappa_0. \tag{4.12}$$

² The implementation of the circle method in its full power would yield the Hardy-Ramanujan-Rademacher type expansions for the multiplicities g_n (cf. [24] where the classical case of $p_{\mathcal{N}}(n)$ is considered). In this paper, we are only concerned with the leading behavior of g_n , and to this end we only need the localization of the integral near the point $z = 1 - \epsilon$. The proof of this property of integral (4.7) is not at all trivial, still this is a rigorous mathematical theorem. Indeed, it needs again the modular properties of the function $f(z)$. It should be also mentioned that there are more general techniques of the asymptotic analysis of the partitions, such as Meinardus theorem (see e.g. [5]; see also [9]). These techniques do not exploit the modular properties of the corresponding generating functions, however, unlike the circle method, they only provide the leading terms of the asymptotics. The Meinardus theorem, as it is stated in [5], is not directly applicable to generating function (4.1).

$\text{Log}[f(\rho)] - n \text{Log } \rho$

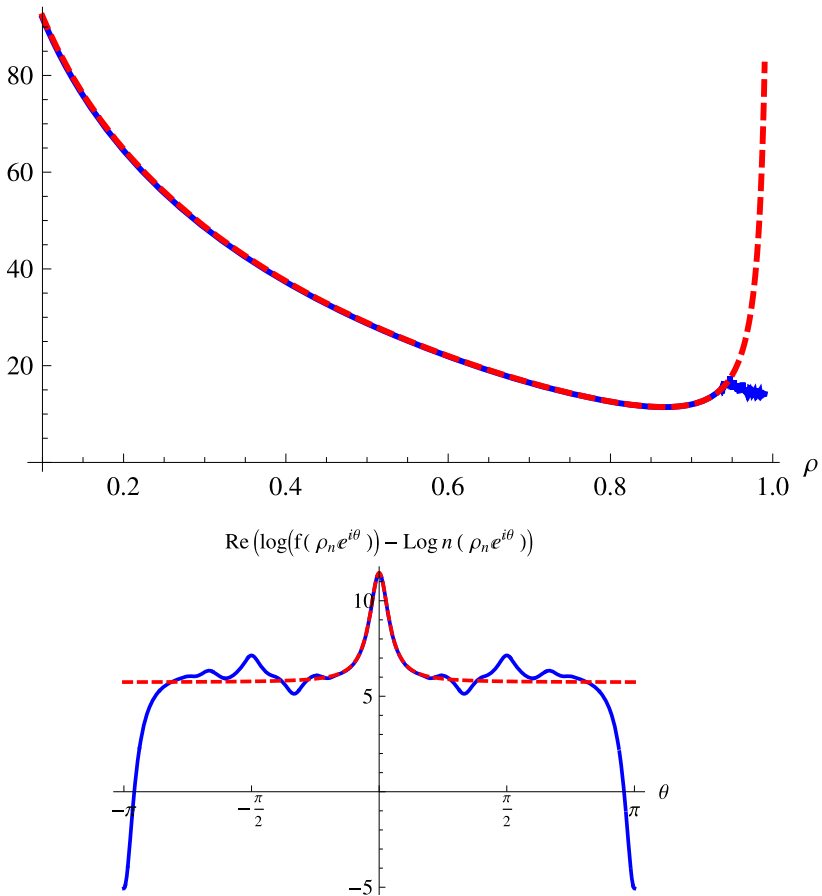


Fig. 2 Case $h > 2$: plots of the logarithm of the generating function in the Cauchy integral (4.7) as a function of the radius ρ (left panel) and of the angular phase θ at the saddle point radius ρ_n given by (4.10) (right panel; only real part shown). The plots show the comparison between the exact expression (4.3) (continuous line, blue) and its asymptotic approximation (4.6) (dashed line, red) at $n = 40$. (color figure online)

Using this specification of δ we re-write estimate (4.11) as

$$\begin{aligned}
 G(z) &= -\frac{\pi}{\sqrt{3}}\sqrt{n - \frac{1}{12}} + \frac{\pi^2}{12\epsilon_n^3}\theta^2 \left(1 + O\left(\epsilon_n^{2\kappa}\right)\right), \\
 &= -\frac{\pi}{\sqrt{3}}\sqrt{n - \frac{1}{12}} + \frac{\pi^2}{12\epsilon_n^3}\theta^2 \left(1 + O\left(n^{-\kappa}\right)\right), \quad 0 < \kappa < \kappa_0, \\
 &n \rightarrow \infty, \quad z \in \mathcal{L}.
 \end{aligned}
 \tag{4.13}$$

In its turn, this estimate yields the following representation for $G(z)$ on \mathcal{L} ,

$$G(z) = -\frac{\pi}{\sqrt{3}}\sqrt{n - \frac{1}{12}} + \frac{\pi^2}{12\epsilon_n^3}t^2(\theta), \tag{4.14}$$

where $t(\theta)$ is a function holomorphic in the neighborhood of the interval $[-\delta, \delta]$ and satisfying the estimates,

$$t(\theta) = \theta (1 + O(n^{-\kappa})), \quad \frac{dt}{d\theta} = 1 + O(n^{-\kappa}), \\ 0 < \kappa < \kappa_0, \quad n \rightarrow \infty, \quad \theta \in [-\delta, \delta]. \tag{4.15}$$

From (4.14) we have that

$$\frac{1}{2\pi i} \int_{\mathcal{L}} \frac{e^{-G(z)}}{z} dz = e^{\frac{\pi}{\sqrt{3}}\sqrt{n - \frac{1}{12}}} \frac{1}{2\pi} \int_{-\delta}^{\delta} e^{-\frac{\pi^2}{12\epsilon_n^3}t^2(\theta)} d\theta. \tag{4.16}$$

At the same time, Eqs. (4.15) allow us to use $t = t(\theta)$ as a new integration variable and transform (4.16) into the asymptotic formula,

$$\frac{1}{2\pi i} \int_{\mathcal{L}} \frac{e^{-G(z)}}{z} dz = e^{\frac{\pi}{\sqrt{3}}\sqrt{n - \frac{1}{12}}} \frac{1}{2\pi} \int_{t(-\delta)}^{t(\delta)} e^{-\frac{\pi^2}{12\epsilon_n^3}t^2} \frac{d\theta}{dt} dt \\ = e^{\frac{\pi}{\sqrt{3}}\sqrt{n - \frac{1}{12}}} \frac{1}{2\pi} \left[\int_{-\delta}^{\delta} e^{-\frac{\pi^2}{12\epsilon_n^3}t^2} dt (1 + O(n^{-\kappa})) + O(n^{-1}) \right], n \rightarrow \infty, \tag{4.17}$$

Assume now that³

$$0 < \kappa < \min \left\{ \frac{1}{4}, \kappa_0 \right\}. \tag{4.18}$$

Then, the integral in the right hand side of (4.17) can be estimated as follows,

$$\frac{1}{2\pi} \int_{-\delta}^{\delta} e^{-\frac{\pi^2}{12\epsilon_n^3}\theta^2} d\theta = \frac{1}{2\pi} \int_{-\infty}^{\infty} e^{-\frac{\pi^2}{12\epsilon_n^3}\theta^2} d\theta - \frac{1}{\pi} \int_{\delta}^{\infty} e^{-\frac{\pi^2}{12\epsilon_n^3}\theta^2} d\theta \\ = \frac{1}{2\pi} \int_{-\infty}^{\infty} e^{-\frac{\pi^2}{12\epsilon_n^3}\theta^2} d\theta - \epsilon_n^{3/2} \frac{1}{\pi} \int_{\delta\epsilon_n^{-3/2}}^{\infty} e^{-\frac{\pi^2}{12}\theta^2} d\theta$$

³ It is worth noticing, that under condition (4.18) the term $O(n^{-1})$ in (4.17) becomes in fact $O(n^{-\infty})$ (exponentially small).

$$\begin{aligned}
 &= \frac{1}{2\pi} \int_{-\infty}^{\infty} e^{-\frac{\pi^2}{12\epsilon_n^3} \theta^2} d\theta + O(n^{-\infty}) \\
 &= \frac{1}{2\pi} \sqrt{\frac{12\epsilon_n^3}{\pi}} + O(n^{-\infty}) = 2^{-3/2} 3^{-1/4} \left(n - \frac{1}{12}\right)^{-3/4} + O(n^{-\infty})
 \end{aligned}
 \tag{4.19}$$

Estimates (4.8), (4.17), and (4.19) yield the following asymptotic formula for the multiplicities g_n of the eigenvalues of the reduced density matrix for $h > 2$.

$$g_n \simeq 2^{-3/2} 3^{-1/4} n^{-3/4} e^{\pi \sqrt{\frac{n}{3}}}, \quad n \rightarrow \infty. \tag{4.20}$$

Turning now to the $h < 2$ case, using (3.14) and remembering that $g_n = 2b_n$, we have

$$\zeta_{\rho_A}(\alpha) = e^{\alpha \left(\frac{1}{6} \ln \frac{k'}{4k^2} - \frac{\pi \tau_0}{6}\right)} f\left(e^{-2\pi \tau_0 \alpha}\right), \tag{4.21}$$

with $f(z)$ defined as in (4.1). As before, using (1.4) we can express $f(z)$ in terms of the Renyi Entropy $S_R(\rho_A, \alpha)$:

$$f(z) = e^{(1-\alpha)S_R(\rho_A, \alpha) - \alpha \left(\frac{1}{6} \ln \frac{k'}{4k^2} - \frac{\pi \tau_0}{6}\right)}, \quad \alpha = -\frac{1}{2\pi \tau_0} \ln z. \tag{4.22}$$

Again, we are interested in the neighbors of the point $z \sim 1$ (see Fig. 3, where we can use the asymptotics derived in [14]):

$$\begin{aligned}
 S_R(\rho_A, \alpha) &= \frac{1}{\alpha(1-\alpha)} \frac{\pi}{12} \frac{I(k)}{I(k')} + \frac{\alpha}{1-\alpha} \frac{1}{6} \ln \frac{k'}{4k^2} + O\left(e^{-\frac{\pi}{\alpha \tau_0}}\right), \\
 \alpha \rightarrow 0, \quad &-\frac{\pi}{2} < \arg \alpha < \frac{\pi}{2}.
 \end{aligned}
 \tag{4.23}$$

Using (4.23) in (4.22) we have (cf. (4.6))

$$f(z) = e^{-\frac{\pi^2}{6 \ln z} - \frac{1}{12} \ln z + O\left(e^{2\pi^2 / \ln z}\right)}, \quad z \rightarrow 1, \quad |z| < 1. \tag{4.24}$$

The figures show that asymptotic agrees with exact analytic expression well. The calculation proceeds exactly as before, where in (4.8) we now have

$$G(z) \equiv \frac{\pi^2}{6 \ln z} + \left(n + \frac{1}{12}\right) \ln z, \tag{4.25}$$

and the saddle point $z = \rho_n$ is:

$$\left. \frac{dG(z)}{dz} \right|_{z=\rho_n} = 0 \quad \Rightarrow \quad \rho_n = e^{-\frac{\pi \sqrt{2}}{\sqrt{12n+1}}} \equiv e^{-\epsilon_n}. \tag{4.26}$$

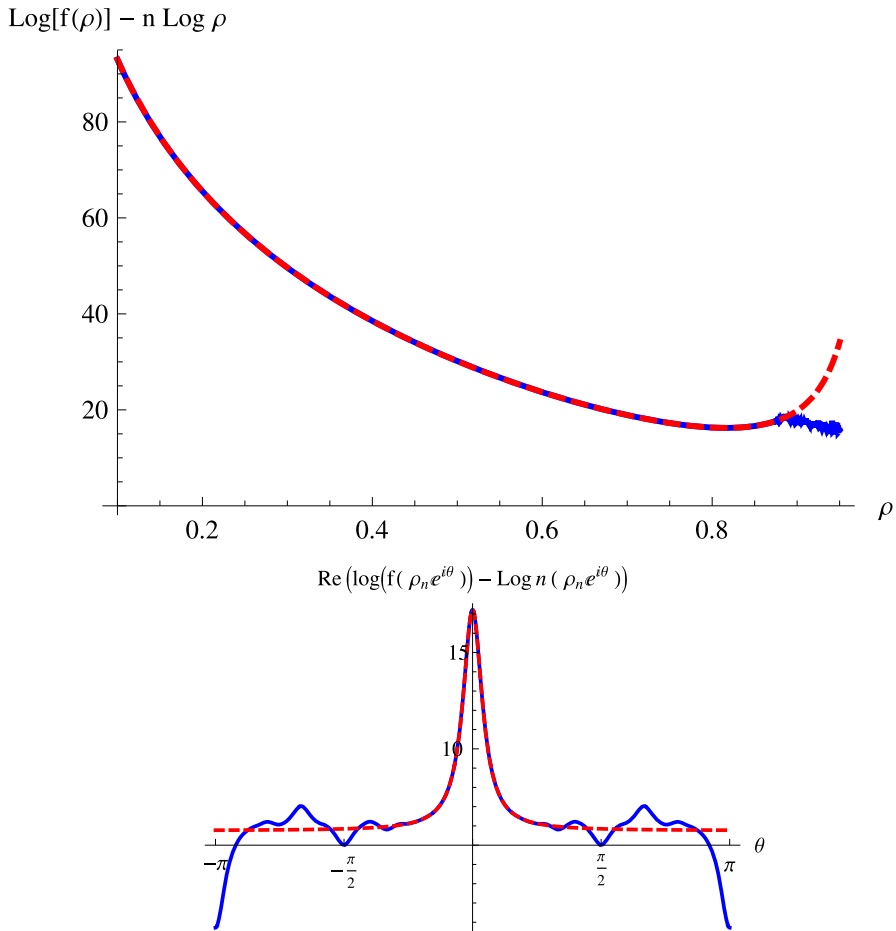


Fig. 3 Case $h < 2$: plots of the logarithm of the generating function in the Cauchy integral (4.7) as a function of the radius ρ (left panel) and of the angular phase θ at the saddle point radius ρ_n given by (4.26) (right panel; only real part shown). The plots show the comparison between the exact expression (4.22) (continuous line, blue) and its asymptotic approximation (4.24) (dashed line, red) at $n = 40$. (color figure online)

Instead of (4.17) we now have,

$$\frac{1}{2\pi i} \int_{\mathcal{L}} \frac{e^{-G(z)}}{z} dz = e^{\pi\sqrt{\frac{2}{3}}(n+\frac{1}{12})} \left[\frac{1}{2\pi} \int_{-\delta}^{\delta} e^{-\frac{\pi^2}{6\epsilon_n^3}\theta^2} d\theta (1 + O(n^{-\kappa})) + O(n^{-1}) \right], \tag{4.27}$$

with δ defined by exactly the same Eq. (4.12). Choosing κ as in (4.18), we can again approximate the integral in the right hand side of (4.27) by a complete Gaussian integral (cf. (4.19)):

$$\begin{aligned} \frac{1}{2\pi} \int_{-\delta}^{\delta} e^{-\frac{\pi^2}{6\epsilon_n^3} \theta^2} d\theta &= \frac{1}{2\pi} \int_{-\infty}^{\infty} e^{-\frac{\pi^2}{6\epsilon_n^3} \theta^2} d\theta + O(n^{-\infty}) \\ &= \frac{1}{2\pi} \sqrt{\frac{6\epsilon_n^3}{\pi}} + O(n^{-\infty}) = 2^{-5/4} 3^{-1/4} \left(n + \frac{1}{12}\right)^{-3/4} \\ &\quad + O(n^{-\infty}) \end{aligned} \tag{4.28}$$

Estimates (4.8), (4.27), and (4.28) yield the following asymptotic formula⁴ for the multiplicities g_n of the eigenvalues of the reduced density matrix for $h < 2$.

$$g_n \simeq 2^{-5/4} 3^{-1/4} n^{-3/4} e^{\pi \sqrt{\frac{2}{3}} n}, \quad n \rightarrow \infty. \tag{4.29}$$

5 Critical lines

As we mentioned in the introduction, there are two critical lines in the phase diagram of the XY model, where the gap closes: at the critical magnetic field $h = 2$ and at the isotropic line $\gamma = 0$ ⁵ and $|h| < 2$. In critical cases ζ function behaves as

$$\zeta_{\rho_A}(\alpha) = \Gamma_{\alpha} \xi^{-\frac{c}{6}(\alpha-1/\alpha)}, \tag{5.1}$$

where Γ_{α} is a non-universal constant and ξ is the relevant length-scale in the considered regime. It was first discovered for isotropic case [18].⁶ The straightforward application of the asymptotic formulae found in [14], shows an agreement with (5.1) close to the critical lines, with the length-scale set by the inverse energy gap Δ , i.e.

$$\xi^{-1} = \Delta = |h - 2|, \quad c = \frac{1}{2}, \quad h \rightarrow 2, \tag{5.2}$$

$$\xi^{-1} = \Delta = |\gamma|, \quad c = 1, \quad \gamma \rightarrow 0, |h| < 2. \tag{5.3}$$

Our Theorems 3.1 and 3.2 for the eigenvalues distributions in the different regimes remain valid arbitrarily close to the critical lines: while the eigenvalues λ_n tend to collapse and vanish with the energy gap, namely

$$\lambda_n \simeq \Delta^{\frac{c}{6}} e^{\frac{\pi^2}{\ln \Delta} n}, \quad \Delta \rightarrow 0, \tag{5.4}$$

their multiplicities do not depend on the parameters γ and h of the model. This agrees with [23]. Our results agree with predictions of conformal field theory.

⁴ It is worth noticing that for the case $h < 2$ the generating function $f(z)$ can be easily transformed to the one satisfying the conditions of the Menardus theorem and hence asymptotics (4.29) can be also obtained by using the Menardus theorem.

⁵ Also known as XX spin chain.

⁶ Renyi entropy was first calculated for XX spin chain.

6 Conclusions

We have calculated the spectrum of the reduced density matrix of the ground state of the anisotropic XY model using the known results on its Renyi entropy [14]. We have confirmed the expectation that, being the model essentially non-interacting, the eigenvalues form a geometric sequence and their multiplicities have simple interpretation in terms of combinatorics and different partitions of integer, see Theorem 3.1 and 3.2. The exact formulae for the eigenvalues have been given in terms of the parameters of the model.

The asymptotic behavior of the multiplicities has been calculated, using the modular properties of Renyi entropy. The leading terms of the asymptotics are given in Eqs. (4.20) and (4.29) for strong and weak magnetic field, respectively. Our results agree with the estimates in [22]. In fact, the log-normal result was achieved by combining (and smearing) the degeneracy with the eigenvalue behavior to give an estimate of the behavior of an effective eigenvalue in a non-integrable system. We think that our method also will work for other spin chains such as XXZ and XYZ . Our results confirm area law [11, 12]

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