# Critical exponents of random XX and XY chains: Exact results via random walks 

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#### Abstract

We study random XY and (dimerized) XX spin- $1 / 2$ quantum spin chains at their quantum phase transition driven by the anisotropy and dimerization, respectively. Using exact expressions for magnetization, correlation functions and energy gap, obtained by the free fermion technique, the critical and off-critical (Griffiths-McCoy) singularities are related to persistence properties of random walks. In this way we determine exactly the decay exponents for surface and bulk transverse and longitudinal correlations, correlation length exponent and dynamical exponent.


Disordered quantum spin chains have gained much interest recently [1]8]. It seems to be established right now that the critical properties in these one-dimensional system are governed by an infinite-disorder fixed-point [9] and the application of a renormalization group (RG) scheme á la Dasgupta and Ma [10] is a powerful tool to determine critical properties and static correlations of these new universality classes, either analytically, if possible, or numerically. Although the underlying renormalization scheme is extremely simple the analytical computations are sometimes tedious [1/[2]. Therefore an alternative route to the exact determination of critical exponents and other quantities of interest is highly desirable, and this is what we are going to present in this letter. In doing so we follow a route on which we already traveled successfully for the random transverse Ising chain 11 13, and here we are going to do one step further studying random XX and XY models with the help of a straightforward and efficient mapping to random walk problems. This mapping is not only a short-cut to the results known from analytical RG calculations, it also gives new exact results in the off-critical region (the Griffiths-phase 14]) and provides a mean to study situations in which the RG procedure must fail, as for instance in the case of correlated disorder 15]. Here we confine ourselves to a concise presentation of the basic ideas including the determination of various exponents for the first time. The technical details of the derivations and further results are deferred to a subsequent publication 16].

The model that we consider is a spin- $1 / 2 \mathrm{XY}$-quantum spin chain with $L$ sites and open boundaries, defined by the Hamiltonian

$$
\begin{equation*}
H=\sum_{l=1}^{L-1}\left(J_{l}^{x} S_{l}^{x} S_{l+1}^{x}+J_{l}^{y} S_{l}^{y} S_{l+1}^{y}\right) \tag{1}
\end{equation*}
$$

where the $S_{l}^{x, y}$ are spin-1/2 operators and the interaction strengths or couplings $J_{l}^{x, y}>0$ are independent random variables modeling quenched disorder. In the case of the
random XY chain one has two independent distributions for the couplings $J^{x}$ and $J^{y}, \rho^{x}$ and $\rho^{y}$, respectively, whereas the random dimerized XX-chain has perfectly isotropic couplings $J_{l}^{x}=J_{l}^{y}=J_{l}$ but two independent probability distributions for the even and odd couplings (i.e. for $J_{2 l}=J_{2 l}^{e}$ and $J_{2 l-1}=J_{2 l-1}^{o}$ ), $\rho^{e}$ and $\rho^{o}$, respectively.

The model (1) has a critical point given by $\left[\ln J^{x}\right]_{\mathrm{av}}=$ $\left[\ln J^{y}\right]_{\mathrm{av}}$ in the XY case and $\left[\ln J^{e}\right]_{\mathrm{av}}=\left[\ln J^{o}\right]_{\mathrm{av}}$ in the XX case (here [...] $]_{\text {av }}$ denotes the disorder average). The distance from the critical point is conveniently measured in the variable

$$
\begin{equation*}
\delta=\frac{\left[\ln J^{x(e)}\right]_{\mathrm{av}}-\left[\ln J^{y(o)}\right]_{\mathrm{av}}}{\operatorname{var}\left[\ln J^{x(e)}\right]+\operatorname{var}\left[\ln J^{y(o)}\right]} \tag{2}
\end{equation*}
$$

where $\operatorname{var}(x)$ is the variance of random variable $x$. At the critical point $(\delta=0)$ spatial correlations decay algebraically, for instance in a finite system of length $L$ with periodic boundary conditions the bulk-correlations decay as

$$
\begin{equation*}
\left[C^{\mu}(L)\right]_{\mathrm{av}}=\left[\langle 0| S_{1}^{\mu} S_{L / 2}^{\mu}|0\rangle\right]_{\mathrm{av}} \sim L^{-\eta^{\mu}} \tag{3}
\end{equation*}
$$

for $\mu=x, y, z,\langle 0|$ denotes the ground state of (1), whereas for a finite system of length $L$ with open boundary conditions the end-to-end correlations decay with a different exponent like

$$
\begin{equation*}
\left[C_{1}^{\mu}(L)\right]_{\mathrm{av}}=\left[\langle 0| S_{1}^{\mu} S_{L}^{\mu}|0\rangle\right]_{\mathrm{av}} \sim L^{-\eta_{1}^{\mu}} \tag{4}
\end{equation*}
$$

Away from the critical point $(\delta \neq 0)$ the infinite system develops long range order. For the $X Y$ model $\lim _{L \rightarrow \infty}\left[C^{\mu}(L)\right]_{\mathrm{av}}=\left(m^{\mu}\right)^{2} \neq 0$, with $m^{x}>0$ for $\delta>0$ and $m^{y}>0$ for $\delta<0$, whereas for the $X X$ model there is dimerization for $\delta \neq 0$ with non-vanishing string order 17]. One can introduce local transverse and longitudinal order parameters $m_{l}^{x, y}$ and $m_{l}^{z}$ also for a finite system (with open boundaries) using the off-diagonal matrix element $\left[m_{l}^{\mu}\right]_{\mathrm{av}}=\left[\langle 1| S_{1}^{\mu}|0\rangle\right]_{\mathrm{av}}$, where $\langle 1|$ is the lowest
excited state with a non-vanishing matrix-element 18]. Analogous to bulk and end-to-end correlations the bulk and surface magnetizations $m_{L / 2}^{\mu}$ and $m_{1}^{\mu}$, respectively, behave differently:

$$
\begin{equation*}
\left[m_{L / 2}^{\mu}\right]_{\mathrm{av}} \sim L^{-x^{\mu}} \quad \text { and } \quad\left[m_{1}^{\mu}\right]_{\mathrm{av}} \sim L^{-x_{1}^{\mu}} \quad(\delta=0) \tag{5}
\end{equation*}
$$

where the critical exponents $x^{\mu}$ and $x_{1}^{\mu}$ fulfill the scaling relation $2 x^{\mu}=\eta^{\mu}$ and $2 x_{1}^{\mu}=\eta_{1}^{\mu}$.

Now we are going to determine the critical exponents introduced above. We use the free-fermion representation of model (11) to derive exact expressions for the local order-parameters whose finite size scaling behavior follows then from a mapping to a random walk problem.

We start with the longitudinal order parameter for the random XX chain (of length $L$ with open boundaries; for convenience we assume from now on that $L$ is a multiple of 4 ), which is given by 16 ]

$$
\begin{align*}
m_{2 l-1}^{z}(X X)=\frac{1}{2}\{1 & +\sum_{k=l}^{L / 2-1} \prod_{j=l}^{k}\left(\frac{J_{2 j-1}}{J_{2 j}}\right)^{2} \\
& \left.+\sum_{k=1}^{l-1} \prod_{j=1}^{k}\left(\frac{J_{2 l-2 j}}{J_{2 l-2 j-1}}\right)^{2}\right\}^{-1} \tag{6}
\end{align*}
$$

for odd sites. Note that the couplings to the left and to the right of the spin that is considered enter this expression differently. For the surface $(l=1)$, where one only has "right" couplings, this expression is similar to an analogous result for the random transverse Ising chain [11] and its scaling properties are related to the survival probability of a random walk with $L / 2$ steps. This is easy to see for the extreme binary distribution, in which $J_{2 j}=1$ and $J_{2 j-1}=\lambda, \lambda^{-1}$ with probability $1 / 2$, taking the limit $\lambda \rightarrow 0$ (i.e. $\lambda^{-1} \rightarrow \infty$ ). Due to the occurance of infinite terms in the sum in the denominator of the r.h.s. of (6) one can easily identify those instances that give a non-vanishing surface magnetization: When $\forall k=1, \ldots, L / 2-1: \prod_{j=1}^{k} J_{2 j-1}<\infty$ the expression on the r.h.s. of ( 6 ) attains a non-vanishing value (typically 1 or, less frequently, some fraction $1 / n$ ), otherwise it is zero. One can represent the disorder configuration $J_{1}, J_{3}, J_{5}, \ldots, J_{L-1}$ as one instance of a random walk with $L / 2-1$ steps by saying that the walker in the $i$-th steps moves downwards if $J_{2 i-1}=\lambda$ and upwards if $J_{2 i-1}=\lambda^{-1}$, as it is sketched in Fig.1. In this way the disorder configuration with non-vanishing surface magnetization $m_{1}^{z}$ are easily identified: they represent surviving walks, i.e. walks that never move into the upper half. Thus $\left[m_{1}^{z}(X X)\right]_{\mathrm{av}}$ scales like the survival probability $P_{\text {surv }}(L / 2)$ of a random walk with $L / 2$ steps that vanishes like $L^{-1 / 2}$, i.e. $\left[m_{1}^{z}(X X)\right]_{\mathrm{av}} \sim L^{-1 / 2}$. Therefore

$$
\begin{equation*}
x_{1}^{z}(X X)=\frac{1}{2} \quad \text { and } \quad \eta_{1}^{z}(X X)=1 \tag{7}
\end{equation*}
$$

Inspecting the expression (6) for the bulk $(l=L / 4)$ one sees that now (again for the extreme binary distribution) a nonvanishing magnetization $m_{L / 2-1}^{z}$ arises


FIG. 1. Sketch of the configuration of odd bonds for a chain of length $L$ that gives a non-vanishing longitudinal magnetization $m_{i}^{z} \sim \mathcal{O}(1)$ for the surface spin, $i=1$, in (a) and the central spin, $i=L / 2-1$, in (b). The example is for the extreme binary distribution with $J_{2 i}=1$. Weak couplings $\left(J_{2 i-1}=\lambda\right)$ correspond to downward steps of the random walk, strong couplings $\left(J_{2 i-1}=\lambda^{-1}\right)$ to upwards steps. The walk in (a) has surviving character, it does not enter the upper half plane. In (b) one can identify two random walks each starting at the central site, $i=L / 2-1$, one to the right and one to the left, and each of them has surviving character.
only if $\forall k=L / 4, \ldots, L / 2-1: \prod_{j=L / 4}^{k} J_{2 j-1}<\infty$ and $\forall k=1, \ldots, L / 4-1: \prod_{j=1}^{k} J_{L / 2-2 j-1}^{-1}<\infty$. We represent the disorder configuration to the right of the central site, $J_{L / 2-1}, J_{L / 2+1}, \ldots, J_{L-1}$, as a random walk with $L / 4$ steps in the way as for the surface spin. The disorder configuration $J_{L / 2-1}, J_{L / 2-3}, \ldots, J_{1}$ to the left is represented as a second (independent) random walk also with $L / 4$ steps, now counting backwards and with the step-direction reversed (i.e. downwards for $J=\lambda^{-1}$ and upwards for $J=\lambda$ ), since now strong bonds on odd sites imply weak coupling of the central spin. For illustration this representation is depicted in Fig. 1. Now, for the bulk magnetization $m_{L / 2-1}^{z}$ to be non-vanishing, both halfs of the coupling configuration have to represent surviving random walks. Thus the probability for a non-vanishing magnetization $m_{L / 2-1}^{z}$ is just the product of two survival probabilities (since both walks are independent), i.e. $\left[m_{L / 2-1}^{z}\right]_{\mathrm{av}} \sim\left\{P_{\text {surv }}(L / 4)\right\}^{2} \sim L^{-1}$ and therefore

$$
\begin{equation*}
x^{z}(X X)=1 \quad \text { and } \quad \eta^{z}(X X)=2 \tag{8}
\end{equation*}
$$

For the XY chain one has the following exact relation [16] for the disorder averaged longitudinal magnetization

$$
\begin{equation*}
\left[m_{l}^{z}(X Y)\right]_{\mathrm{av}}=\left[\left\{m_{l}^{z}(X X)\right\}^{1 / 2}\right]_{\mathrm{av}}^{2} \tag{9}
\end{equation*}
$$

which yields immediately the surface and bulk exponents
for longitudinal order parameter and correlations: Since $\left\{m_{1}^{z}(X X)\right\}^{1 / 2}$ has a non-vanishing value if and only if $m_{1}^{z}(X X)$ is non-vanishing, one obtains $\left[m_{1}^{z}(X Y)\right]_{\mathrm{av}} \sim$ $\left\{P_{\text {surv }}(L / 4)\right\}^{2} \sim L^{-1}$, which means

$$
\begin{equation*}
x_{1}^{z}(X Y)=1 \quad \text { and } \quad \eta_{1}^{z}(X Y)=2 \tag{10}
\end{equation*}
$$

and further, $\left[m_{L / 2-1}^{z}(X Y)\right]_{\mathrm{av}} \sim\left\{P_{\text {surv }}(L / 4)\right\}^{4} \sim L^{-2}$, which means

$$
\begin{equation*}
x^{z}(X Y)=2 \quad \text { and } \quad \eta^{z}(X Y)=4 \tag{11}
\end{equation*}
$$

For the transverse surface order parameter $m_{1}^{x}$ one has [16] the exact formula (valid for the XX and XY case)

$$
\begin{equation*}
m_{1}^{x}=\frac{1}{2}\left[1+\sum_{l=1}^{L / 2-1} \prod_{j=1}^{l}\left(\frac{J_{2 j-1}^{y(o)}}{J_{2 j}^{x(e)}}\right)^{2}\right]^{-1 / 2} \tag{12}
\end{equation*}
$$

which is similar to (6) for $m_{1}^{z}(X X)$, apart from the power $1 / 2$ on the r.h.s., which again is only non-vanishing for disorder configurations that represent surviving walks á la Fig.1a. This then yields without any effort

$$
\begin{equation*}
x_{1}^{x}(X X, X Y)=1 \quad \text { and } \quad \eta_{1}^{x}(X X, X Y)=2 \tag{13}
\end{equation*}
$$

For the transverse bulk order parameter in the XY chain we use the fact that the model can be mapped onto two transverse Ising models (TIM), with uncorrelated disorder in both chains [2,16]. Through this mapping one obtains for the transverse correlation function $C_{2 i, 2 i+2 r}^{x}=\langle 0| S_{2 i}^{x} S_{2 i+2 r}^{x}|0\rangle$ the following identity 116]

$$
\begin{aligned}
{\left[C_{2 i, 2 i+2 r}^{x}(X Y)\right]_{\mathrm{av}}=} & 4\left[C_{i, i+r}^{x}\left(T I M_{\mathrm{free}}\right)\right]_{\mathrm{av}} \\
& \cdot\left[C_{i, i+r}^{x}\left(T I M_{\mathrm{fixed}}\right)\right]_{\mathrm{av}} \sim r^{-2 \eta^{x}(T I M)},
\end{aligned}
$$

where fixed and free indicated the boundary conditions. Since the correlation function exponent is known exactly [1] to be $\eta^{x}(T I M)=(3-\sqrt{5}) / 2$ we have:

$$
\begin{equation*}
x^{x}(X Y)=(3-\sqrt{5}) / 2 \quad \text { and } \quad \eta^{x}(X Y)=3-\sqrt{5} \tag{15}
\end{equation*}
$$

For the XX chain the two transverse Ising chains have perfectly correlated disorder, which implies that the disorder averaged transverse correlations do not factorize into two independent averages as in (14). Therefore, for the transverse order parameter exponent in the XX case we have to use a different route: The first important observation is that the transverse bulk order parameter $m_{L / 2}^{x}=\langle 1| S_{L / 2}^{x}|0\rangle$ attains its maximum value $1 / 2$ if the central spin is decoupled from the rest of the system, i.e. when $J_{L / 2-1}=J_{L / 2}=0$. More generally we expect that $m_{L / 2}^{x} \sim \mathcal{O}(1)$ when it is weakly coupled to the rest of the system. "Weakly coupled" in the case of the extreme binary distribution means that the bond configuration to the left and to the right of the central spin represent both surviving random walks, as exemplified in Fig.1b (this is actually equivalent to
the (exact) condition for the longitudinal order parameter $m_{L / 2}^{z}(X X)$ to be non-vanishing). This correspondence implies $\left[m_{L / 2}^{x}(X X)\right]_{\mathrm{av}} \sim\left\{P_{\text {surv }}(L / 4)\right\}^{2} \sim L^{-1}$ from which one obtains

$$
\begin{equation*}
x^{x}(X X)=1 \quad \text { and } \quad \eta^{x}(X X)=2 \tag{16}
\end{equation*}
$$

We verified the strong correlation between weak coupling and non-vanishing transverse order parameter numerically in the following way: We considered a chain with $L+1$ sites and the couplings at both sides of the central spin were taken randomly in the form of surviving walk character, where we used the binary distribution with $\lambda=0.1$. For such small value of $\lambda$ the surface orderparameter averaged over the surviving walk (sw) configurations $\left[m_{1}^{x}\right]_{\text {sw }}$ was very close to the maximal value of $1 / 2$. Then we calculated numerically the order-parameter at the central spin and its average value over surviving walk configurations $\left[m_{L / 2}^{x}\right]_{\mathrm{sw}}$ as given in Table I.

| L | $2\left[m_{1}^{x}\right]_{\mathrm{sw}}$ | $2\left[m_{L / 2}^{x}\right]_{\mathrm{sw}}$ |
| :---: | :---: | :---: |
| 32 | 0.994 | 0.764 |
| 64 | 0.991 | 0.682 |
| 128 | 0.991 | 0.647 |
| 256 | 0.991 | 0.577 |

TABLE I: Surface and bulk transverse order-parameters averaged over 50000 surviving walk coupling configurations for the binary distribution $(\lambda=0.1)$.

As seen in the Table the averaged surface orderparameter stays constant for large values of $L$, whereas the bulk order-parameter decreases very slowly, actually slower than any power. The data can be nicely fitted by $\left[m_{L / 2}^{x}\right]_{\mathrm{sw}} \sim(\ln L)^{-1 / 2}$. Thus we conclude that the numerical results confirm the exponents given in (16), however there are strong logarithmic corrections, which imply for the average transverse correlations

$$
\begin{equation*}
\left[C^{x}(r)\right]_{\mathrm{av}} \sim r^{-2} \ln ^{-1}(r) \quad \mathrm{XX}-\text { model } \tag{17}
\end{equation*}
$$

These strong logarithmic corrections render the numerical calculation of critical exponents very difficult 17,19. In earlier numerical work using smaller finite systems disorder dependent exponents were reported 19. We believe that these numerical results can be interpreted as effective, size-dependent exponents and the asymptotic critical behavior is indeed described by Eq. (17).

Away from the critical point the correlation length exponent $\nu$ can be determined by the scaling behavior of the longitudinal surface magnetization, i.e. (6) with $l=1$, which can be inferred from the survival properties of a, now biased, random walk: $\left[m_{1}^{x, y}(\delta, L)\right]_{\mathrm{av}} \sim P_{\text {surv }}(\delta, L / 2)$. A non-vanishing distance $\delta$, see (2), from the critical point means that the disorder configurations can be represented by random walk that have a drift either towards $(\delta>0)$ or away from $(\delta<0)$ an absorbing wall (take for instance the extreme binary distribution, in which weak
bonds $\lambda$ occur with a probability $(1-\delta) / 2$ and strong bonds $\lambda^{-1}$ occur with a probability $(1+\delta) / 2$ and compare with Fig.1a). Recalling the asymptotic properties of the survival probability of random walks [11] one gets for $\delta>0 P_{\text {surv }}(\delta>0, L) \sim \exp (-L / \xi)$ with $\xi \sim \delta^{-2}$, where the characteristic length scale $\xi$ of surviving walks corresponds to the average correlation length of the XX and XY chains:

$$
\begin{equation*}
[\xi]_{\mathrm{av}} \sim \delta^{-\nu} \quad \text { with } \quad \nu=2 \tag{18}
\end{equation*}
$$

For $\delta<0$ the drift away from the adsorbing wall yields a finite survival probability even in the infinite system, $P_{\text {surv }}(\delta<0, L / 2) \propto \delta$, which implies that $\left[m_{1}^{\mu}\right]_{\mathrm{av}} \propto|\delta|^{-\beta_{1}^{\mu}}$, with $\beta_{1}^{x, y}(X X, X Y)=\beta_{1}^{z}(X X)=1$ and $\beta_{1}^{z}(X Y)=2$. Since $x_{1}^{\mu}=\beta_{1}^{\mu} / \nu$ and $\nu=2$ from (18) one reconfirms the results about the surface magnetization exponents in (7), (10) and (13)

The typical correlation length, $\xi_{\text {typ }}$ can be inferred from the scaling behavior of the typical surface magnetization $\ln m_{1} \sim \sum_{j}\left\{\ln \left(J_{2 j-1}^{y(o)}\right)-\ln \left(J_{2 j}^{x(e)}\right)\right\} \propto \delta L$, which gives

$$
\begin{equation*}
[\xi]_{\text {typ }} \sim \delta^{-\nu_{\text {typ }}} \quad \text { with } \quad \nu_{\text {typ }}=1 \tag{19}
\end{equation*}
$$

The critical and off-critical scaling behavior of the low energy exciations and dynamical correlations can be deduced from the formula for the gap $\epsilon$ 16]

$$
\begin{equation*}
\epsilon(L)=m_{1}^{x} m_{L-1}^{x} J_{L-1}^{y} \prod_{j=1}^{L / 2-1} \frac{J_{2 j-1}^{y(o)}}{J_{2 j}^{x(e)}}, \tag{20}
\end{equation*}
$$

which is analogous to a corresponding formula for the gap in the transverse Ising chain 11]. At the critical point one observes that $\ln \epsilon$ is a sum of $L$ independently distributed random variables with zero mean (since $\delta=0$ ), for which the central limit theorem applies. Therefore the probability distribution of gaps obeys $P(\ln \epsilon) \sim L^{-1 / 2} \tilde{p}\left(\ln \epsilon / L^{-1 / 2}\right)$ and one uses scaling arguments as in 20] to deduce the asymptotic (imaginary) time dependence of the spin-spin autocorrelation function $G_{l}^{\mu}(\tau)=\left[\langle 0| S_{l}^{\mu}(\tau) S_{l}^{\mu}(0)|0\rangle\right]_{\mathrm{av}}$.

$$
\begin{equation*}
G_{a}^{\mu}(\tau) \sim(\ln \tau)^{-\eta_{a}^{\mu}} \tag{21}
\end{equation*}
$$

for the surface ( $a=1$ ) and bulk ( $a=\mathrm{bulk}$ ), respectively, with the critical exponents $\eta_{a}^{\mu}$ as given above.

Away from the critical point in the Griffiths-phase 14] the gap distribution has still an algebraic tail $P(\epsilon) \sim$ $\epsilon^{-1+1 / z^{\prime}(\delta)}$, with a dynamical exponent $z^{\prime}(\delta)$ that varies continuously with the distance from the critical point $\delta$ and is given by the exact (implicite) formula 12

$$
\begin{equation*}
\left[\left(\frac{J^{x(e)}}{J^{y(o)}}\right)^{1 / z^{\prime}(\delta)}\right]_{\mathrm{av}}=1 \tag{22}
\end{equation*}
$$

The dynamical exponent $z^{\prime}(\delta)$ parameterizes all GriffithsMcCoy singularities occurring in the Griffiths-phase, e.g. the spin-spin autocorrelations decay algebraically as

$$
\begin{equation*}
G_{l}(\tau) \sim \tau^{-1 / z^{\prime}(\delta)} \tag{23}
\end{equation*}
$$

which gives for the susceptibility $\chi^{\mu} \sim T^{-1+1 / z^{\prime}(\delta)}$ diverging for $T \rightarrow 0(T=$ temperature $)$ if $z^{\prime}(\delta)>1$.

To summarize we have shown how to obtain a complete description of the critical and off-critical singularities of random XX and XY chains with simple random walk arguments using exact formulas arising from the free fermion description of these quantum spin models. All results for the critical exponents are therefore exact. One should note that for the transverse bulk order parameter exponent for the XY model we referred to a result for the transverse Ising model obtained by a RG calculation [1] and for the same exponent of the XX model we showed the existence of strong logarithmic corrections.
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