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Exact Haldane mapping for all \( S \) and super universality in spin chains

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Abstract – The low-energy dynamics of the anti-ferromagnetic Heisenberg spin \( S \) chain in the semiclassical limit, \( S \to \infty \) is known to map onto the \( O(3) \) nonlinear \( \sigma \)-model with a \( \theta \) term in \( 1+1 \) dimension. Guided by the underlying dual symmetry of the spin chain, as well as by the recently established topological significance of “dangling edge spins”, we report an exact mapping onto the \( O(3) \) model that avoids the conventional large-\( S \) approximation altogether. Our new methodology demonstrates all the super universal features of the \( \theta \)-angle concept that previously arose in the theory of the quantum Hall effect. It explains why Haldane’s original ideas remarkably yield the correct answer in spite of the fundamental complications that generally exist in the idea of semiclassical expansions.

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In 1983, Haldane proposed that the low-energy dynamics of the anti-ferromagnetic Heisenberg spin \( S \) chain can be taken from the \( O(3) \) nonlinear \( \sigma \)-model (NLSM) in \( 1+1 \) dimension and in the presence of the \( \theta \) term [1]. At least within the limitations of a large-\( S \) approximation the parameter \( \theta \) was found to take on the values 0 or \( \pi \) only, depending on whether \( S \) is integral or half-integral, respectively. Since the standard \( O(3) \) model with \( \theta = 0 \) is known to display a mass gap [2], Haldane concluded that the integral spin chain is always gapped. This is unlike the \( S = 1/2 \) chain, for example, which from the Bethe ansatz solution is known to display gapless excitations [3]. Based on numerical work it is now generally accepted that the uniform integral spin chain is always gapped, whereas the half-integral spin chain is generally gapless.

It has recently been pointed out, however, that a universal topological feature of the \( \theta \)-angle concept has historically been overlooked [4]. The \( \theta \)-vacuum quite generally displays massless chiral edge excitations [5] that have important consequences for the theory on the strong-coupling side. Within the Grassmannian \( U(M+N)/U(M) \times U(N) \) nonlinear sigma model approach to localization and interaction phenomena, for example, it was shown that the massless edge excitations are directly related to the existence of robust topological quantum numbers that explain the stability and precision of the quantum Hall effect [6]. This has led to the idea of super universality of quantum Hall physics that, unlike the common belief in the field, is independent of the details of the theory such as the number of field components and in particular, the replica limit \( M, N \to 0 \) [6].

Based on the Haldane mapping it is next natural to expect that the super-universality statement can be extended to also include quantum spin liquids. Indeed, renormalization group studies have clearly indicated that the dimerised \( S = 1/2 \) spin chain displays all the basic features of the quantum Hall effect [4] in much the same manner as what has recently been observed, for example, in the exactly solvable large-\( N \) limit of the \( CP^{N-1} \) model [7]. The dynamics of the dangling edge spins of the chain thereby plays a role that is in many ways the same as that of the massless chiral edge excitations of the \( \theta \)-vacuum. Unfortunately, a generalized Haldane mapping that includes the dangling spins at the edges of the spin \( S \) chain has so far been obtained only for the semiclassical \( S = \infty \) limit [4]. Given the extensive literature on the subject of quantum spin chains, it is somewhat surprising to know that not a single attempt has been reported that
would in principle resolve this longstanding drawback and extend the Haldane mapping to include finite values of $S$.

One of the main objectives of the present letter is to show that the statement of super universality of quantum spin liquids cannot in general be established by using the semiclassical large-$S$ idea alone. To illustrate the problems we first recall the Hamiltonian of the spin chain

$$\mathcal{H} = \frac{J}{S} \sum_I (S_{I1} \cdot S_{I2} + \kappa S_{I2} \cdot S_{I+1}).$$

(1)

Here, $J > 0$ favors the anti-ferromagnetic (Néel) ordering, the integer $I$ is the dimer index and the subscripts 1 and 2 denote the two sites within each dimer. The variable $\kappa$ with $0 < \kappa < \infty$ is the dimensionless nearest-neighbor coupling between the dimers. The path integral representation involves $O(3)$ vectors $\mathbf{n}(t)$ with $\mathbf{n}^2 = 1$ and $t$ denoting the imaginary time. The action is [8]

$$S = iS \sum_I \Omega[\mathbf{n}_{I1}] + \Omega[\mathbf{n}_{I2}]$$

$$+ SJ \sum_I \oint (\mathbf{n}_{I1} \cdot \mathbf{n}_{I2} + \kappa \mathbf{n}_{I2} \cdot \mathbf{n}_{I+1}).$$

(2)

The quantity $\Omega[\mathbf{n}_{I0}]$ is the solid-angle term associated with each lattice site $I0$, and $\oint = \int_0^\beta dt$. To extract the low-energy dynamics of the spin chain from eq. (2) several basic assumptions are necessary. The historical procedure was based on the change of variables [9]

$$\mathbf{m}_I = \frac{1}{2} (\mathbf{n}_{I1} - \mathbf{n}_{I2}), \quad \lambda_I = \frac{1}{2} (\mathbf{n}_{I1} + \mathbf{n}_{I2}),$$

(3)

where $\mathbf{m}_I = \mathbf{m}_I/|\mathbf{m}_I|$, which is a measure of the Néel ordering, is taken as the soft mode in the problem that should be retained. The field variable $\lambda_I$, on the other hand, is taken as the hard mode that should be integrated out. The simplest way to do this is by using semiclassical approximations. Assuming $S \to \infty$ then eq. (2) can be evaluated at the saddle point which in the long-wavelength limit (slowly varying $\mathbf{m}_I$) yields the NLSM in the presence of the $\theta$-angle [4,9]. The complications, however, occur in the computation of the $1/S$ corrections that all diverge as $\beta \to \infty$. The origin of the divergences is easily understood from the following exact expression for the solid-angle term of each dimer [10]:

$$\Omega[\mathbf{n}_{I1}] + \Omega[\mathbf{n}_{I2}] = 2 \oint \lambda_I \cdot \mathbf{m}_I \times \partial_t \mathbf{m}_I.$$

(4)

The action of the “hard” modes has therefore no time derivatives and this, in turn, implies that coincident operators $\lambda_I(t)$ have a divergent expectation value. What this simple example is telling us is that semiclassical expansions are quite generally plagued by ambiguities that are inherent in the bosonic path integral representation of eq. (2). These ambiguities, as we shall see later on in this letter, are the main reason why the general theory of the Heisenberg anti-ferromagnet has not advanced beyond the naive $S = \infty$ saddle point limit.

**Three-spin problem.** As the most important step next we introduce a novel mapping onto the NLSM that avoids the use of semiclassical approximations altogether. This is accomplished if, instead of eq. (3), we pursue an effective action of the spin chain as defined by decimation, i.e. by eliminating the spins on one sublattice (say, $\mathbf{n}_{I1}$) while retaining those on the other ($\mathbf{n}_{I2}$). The primary focus will be on open spin chains containing $2N + 1$ sites since they display both the subtleties of “dangling edge spins” and “dual invariance.” By dual invariance we mean that the dimerised spin system described by eqs. (1) and (2) is invariant under the interchange between the weak and strong bonds

$$J \to J\kappa, \quad \kappa \to \frac{1}{\kappa}.$$

(5)

For simplicity we first consider the three-spin system as sketched in fig. 1 which is the simplest possible system with edges that is manifestly self-dual. It is advantageous to express the $O(3)$ vector field $\mathbf{n}$ in terms of a Grassmannian field variable $Q = \mathbf{n} \cdot \tau \in SU(2)/U(1)$, with $\tau$ denoting the Pauli matrices. This leads to the three-spin action

$$S[Q_1, Q_2, Q_3] = iS \left\{ \Omega[Q_1] + \Omega[Q_2] + \Omega[Q_3] \right\}$$

$$+ \frac{1}{2} SJ \oint \text{tr}(Q_1 + \kappa Q_3) Q_2.$$

(6)

We can always write $Q = T^{-1} T$, with $T \in SU(2)$ such that the solid-angle term can be expressed explicitly according to $i\Omega[Q] = i\Omega[T] = \oint \text{tr} \partial_t T^{-1} T^{-1} \partial_t T$ [4]. The decimation of the three-spin system is defined as follows:

$$e^{-S_{eff}[Q_1, Q_2]} = \int \mathcal{D}[Q_2] e^{-S[Q_1, Q_2, Q_3]}.$$

(7)

To evaluate $S_{eff}$ we notice that the matrix $Q_1 + \kappa Q_3$ in eq. (6) is Hermitian and traceless. Therefore, let us write

$$J (Q_1 + \kappa Q_3) = BV^{-1} T V,$$

(8)

$$B = J (1 + \kappa) \sqrt{1 - \frac{\kappa}{2(1 + \kappa)^2}} \text{tr} Q Q,$$

(9)

where $V \in SU(2)$ and $\delta T = Q_3 - Q_1$ which is taken as a small quantity. Similarly, we expand the matrix $V$ in powers of the differential $\delta T = T_3 - T_1$,

$$V = T_1 + \frac{\kappa}{1 + \kappa} \delta T + O(\delta^2).$$

(10)

Based on eqs. (8)–(10) we obtain $S_{eff}$ order by order in a derivative expansion. To see this we replace

Fig. 1: The 3-spin problem with relative bond strengths 1 and $\kappa$. 

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\( \theta_{\binom{1}{2}} = \theta_{\binom{1}{2}} + 2\pi k(\binom{1}{2}) \),

(20)

where \(-\pi < \theta_{\binom{1}{2}} \leq \pi\) denotes the fractional piece and \(k(\binom{1}{2}) = 0, 1, \ldots, 2N\) the integral piece of \(\theta(\binom{1}{2})\). Notice that under the transformation of eq. (5), the NLSM parameters are replaced by

\[
\theta_{\binom{1}{2}}(1/\binom{1}{2}) = -\theta_{\binom{1}{2}}(k(1/\binom{1}{2})), \quad k(1/\binom{1}{2}) = S - k(\binom{1}{2}),
\]

(21)

Equation (20) therefore permits a splitting of eq. (17) into a “bulk” part \(S_{\binom{1}{2}} [Q] \) and an “edge” part \(S_{\binom{1}{2}} [T] \) that are decoupled under the dual transformation. Specifically, we rewrite

\[
S_{\binom{1}{2}}[T] = \frac{k}{2} \Omega[T(0)] + i \left( S - \frac{k}{2} \right) \Omega[T(L)].
\]

(24)

These final expressions are the principal results of this letter. Notice that eq. (24), is a fundamental statement made on the “edge” of the-th-vacuum that has traditionally gone unnoticed [4,6]. Following eqs. (12) and (14) it is the critical action of “dangling edge spins” and should therefore be regarded as an integral part of the low-energy dynamics of the spin chain [11]. To understand this aspect of the problem we first consider the class of topological field configurations \(Q_0 \) for which \(C[Q_0] \) is strictly an integer. As is well known, this class sets the stage for the “instanton picture” of the-th-angle and is geometrically defined by identifying the edge as a single point, i.e. \(Q_0 \) is a constant matrix along the edge of the system, say \(Q_0 = \tau_z \) [6]. The action for the edge equation (24) is now a trivial phase factor and we immediately recognize \(S_{\binom{1}{2}}[T] = S_{\binom{1}{2}}[Q_0] \) as the theory of the “bulk” of the system that only depends on \(\theta(\binom{1}{2}) \) modulo \(2\pi\).

To incorporate the edge excitations we write \(Q = U^{-1}Q_0 U \). Here \(U \in SU(2) \) represents the “fluctuations” about the boundary condition \(Q_0 = \tau_z \) that carry a fractional topological charge. Discarding unimportant phase factors eq. (24) now reads \(S_{\binom{1}{2}}[T] = S_{\binom{1}{2}}[U] \) indicating that the matrix \(U \) is the basic field variable for the “edge.” Of primary interest is the effective action for “edge” excitations \(S_{\binom{1}{2}}[U] \) which is given by [4,6]

\[
e^{-S_{\binom{1}{2}}[U]} = e^{-S_{\binom{1}{2}}[U]} \int_{\partial V} D[Q_0] e^{-S_{\binom{1}{2}}[U^{-1}Q_0 U]}.
\]

(25)

Here, \(\partial V \) reminds us of the boundary condition \(Q_0 = \tau_z \). Next, we make use of the fact that the two-dimensional

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As explained later in the text the same expressions are obtained in the naive \(S = \infty \) approximation [4,9].
theory of eq. (23) develops a mass gap when \( \theta_B(\kappa) \approx 0 \) or \( \theta(\kappa) \approx 2\pi k(\nu) \) [2]. A finite mass gap in the "bulk" means that the functional integral of eq. (25) is insensitive to changes in the boundary conditions and, hence, \( \tilde{S}_E[U] \equiv S_E[U] \) except for \( U \)-independent terms. We therefore conclude that the dimerised spin chain with varying values of \( \kappa \) displays 2\( S+1 \) different topological phases, labelled by the integer \( k(\kappa) \), where the low-energy excitations are solely those of the "dangling edge spins" with quantum numbers \( k(\kappa)/2 \) and \( S-k(\kappa)/2 \), respectively. These phases must in general be separated by quantum phase transitions (or a vanishing mass gap in the bulk) occurring at intermediate values of \( \kappa \) where \( \theta_B(\kappa) \) makes a "jump" from \( +\pi \) to \( -\pi \). Notice that gapless excitations are in general necessary in order for the spin chain to be able to transport a spin-\( \frac{1}{2} \) quantum over macroscopic distances from one edge to the other. Moreover, in complete analogy with the "electrodynamics picture" of the \( \theta \)-angle [12] one may interpret the "jump" in \( \theta_B(\kappa) \) in terms of the creation of "Coleman charges" that move to the opposite edges so as to maximally shield the "background electric field" \( \theta(\kappa) \) (see footnote \(^2\)).

In summary, emerging from eqs. (22)–(25) are precisely the super universal features of the \( \theta \)-vacuum concept that have previously been discovered in the context of the quantum Hall effect [6]. The statement of super universality becomes all the more transparent when, for example, as well as to the Heisenberg antiferromagnet in higher dimensions [14].

Semiclassical theory. – Next, it is of interest to know what our explicit and exact results teach us about the problem of semiclassical expansions that to date has spanned the subject. For this purpose we address the single-spin problem of eq. (12) and employ the Schwinger boson representation \( Q^z_{\alpha} = \delta_{\alpha\sigma} - 2z_{\alpha} z^*_{\sigma} \), with \( z^* \cdot z = 1 \). Fixing the \( U(1) \) gauge such that \( z_1 = z^*_1 = \sqrt{1-z^2_{22}} \), the \( Q_2 \) matrix field can be written as

\[
Q_2 = \begin{pmatrix}
1 - 2z^2_{22} & 2z_{22}\sqrt{1-z^2_{22}} \\
2z^*_{22}\sqrt{1-z^2_{22}} & -1 + 2z^2_{22}
\end{pmatrix}.
\]

The single-spin action of eq. (12) becomes simply

\[
S_0 = 2S \int \left\{ z^*_2(t)\partial_t z_2(t) + B_0 z^*_2(t-\epsilon)z_2(t) \right\}. \tag{28}
\]

For reasons to be explained shortly, we have introduced an infinitesimal time-splitting quantity \( \epsilon > 0 \). In the definition of \( S_0 \). In the large-\( S \) limit the propagator becomes

\[
\langle z^*_2(0)z_2(t) \rangle = \frac{1}{2S} \delta(t-\epsilon)e^{-B_0\epsilon}. \tag{29}
\]

It can be shown that eqs. (27)–(29) are completely equivalent to the Holstein-Primakoff representation of the single spin. The effect of our introduction of the time-splitting quantity \( \epsilon \) in eqs. (28) and (29) is to render the expectation value of coincident operators identically equal to zero. The remarkable conclusion that one can draw from all this is that the exact correlations of eq. (14) are obtained in the limit \( S \to \infty \) and the corrections are zero to all orders in \( 1/S \) [13].

The subtle but crucial feature of coincident operators is generally lost in eqs. (2) and (12), a drawback of the bosonic path integral that explains why the traditional saddle point or large-\( S \) technique is complicated. On the other hand, the present theory fundamentally resolves these ambiguities since our exact and semiclassical results can be compared directly with the knowledge obtained from different sources, notably the Hamiltonian formalism. Finally, our methodology is not specifically designed for \( SU(2) \) spins alone. It can also be applied to the \( SU(n) \) case, for example, as well as to the Heisenberg antiferromagnet in higher dimensions [14].

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REFERENCES

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