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Exact Dynamical Structure Factor of the Degenerate Haldane-Shastry Model

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The dynamical structure factor $S(q,\omega)$ of the K-component ($K = 2$, 3, 4) spin chain with a $1/r^2$ interaction is derived exactly at zero temperature for the arbitrary size of the system. The result is interpreted in terms of a free quasiparticle picture which is a generalization of the spinon picture in the SU(2) case. The excited states consist of K quasiparticles each of which is characterized by a set of $K - 1$ quantum numbers. Divergent singularities of $S(q,\omega)$ at the spectral edges are derived analytically. The analytic result is checked numerically for finite systems.

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Recently much interest has been focused on magnetic systems with orbital degeneracy [1–6]. In the case of twofold orbital degeneracy, the total degeneracy per site becomes $4(= 2 \times 2)$, and the simplest model to realize this situation in one dimension is the spin chain with SU(4) symmetry. The static property of the SU(4) spin chain has been studied mainly by numerical methods. It has been reported that the spin correlation has a period of four unit cells, and that the asymptotic decay has a power-law exponent different from unity [3,4]. Such an exponent has also been derived by use of conformal field theory [7]. In view of this situation, one can naturally ask how the dynamical property depends on the number of internal degrees of freedom. Experimental investigations of orbitally degenerate quasi-one-dimensional magnetic compounds such as NaV_2O_5 [8] are being performed with increasing accuracy. Hence it is useful to clarify the difference from systems without orbital degeneracy not only for static properties but also for dynamic ones.

In this Letter, we derive exact an analytic formula for the dynamical structure factor $S(q,\omega)$ of the SU(K) spin chain at zero temperature, and provide intuitive interpretation of the result in terms of quasiparticles obeying fractional statistics. We take the exchange interaction J_{ij} decaying as an inverse square of the distance: $J_{ij} = J[(N/\pi)\sin\pi(i-j)/N]^{-2}$, where N is the number of lattice sites with unit spacing and $J > 0$. The model is given by [9,10]

$$H_{\mathrm{HS}} = \frac{1}{2}\sum_{1\le i<j\le N} J_{ij}P_{ij}\,, \tag{1}$$

where P_{ij} is the exchange (or permutation) operator. It can be written in the form

$$P_{ij} = \sum_{\delta,\gamma=1}^{K} X_i^{\delta\gamma}X_j^{\gamma\delta}, \tag{2}$$

where $X_i^{\gamma\delta}$ changes the spin state δ to γ at site i. In the particular case of SU(2), P_{ij} is reduced to the spin exchange $2\vec{S}_i\cdot\vec{S}_j + 1/2$. This model is a generalization of the Haldane-Shastry (HS) model [11,12] for the SU(2) chain, and hence is called the SU(K) HS model in the following.

In the original HS model, the spinons form an ideal spin 1/2 "semion" gas [13], obeying the fractional exclusion statistics [14]. The dynamical structure factor $S(q,\omega)$ of the SU(2) HS model has a remarkably simple structure in terms of the spinon picture: only two spinons contribute to $S(q,\omega)$ [15]. Since the semionic statistics is applicable only to the case of SU(2), one has to take more general fractional statistics in order to apply a quasiparticle description.

To derive the exact formula for the dynamical structure factor, we use the U(K) spin Calogero-Sutherland (CS) model [10] as an auxiliary. The Hamiltonian of the U(K) spin CS model is given by

$$H_{\mathrm{spinCS}} = -\frac{1}{2}\sum_{i=1}^{N}\frac{\partial^2}{\partial x_i^2} + \left(\frac{\pi}{L}\right)^2\sum_{1\le i<j\le N}\frac{\beta(\beta+P_{ij})}{\sin^2\frac{\pi}{L}(x_i-x_j)}\,, \tag{3}$$

where $\beta > 0$ is the coupling parameter and L is the size of the system. This continuous model is more tractable than the SU(K) HS model, because the eigenfunctions of the model have been explicitly constructed [16,17]. We take the strong coupling limit $\beta\to\infty$ of the U(K) spin CS model. Then particles crystallize with the lattice parameter L/N which is taken as the unit of length. We are then left with the center-of-mass motion, the lattice vibration and the dynamics of the internal degrees of freedom which is called the "color." The color dynamics is equivalent to the dynamics of the SU(K) HS model. The freezing trick described above was first introduced by Polychronakos [18], and has been applied to thermodynamics of lattice models [19–21]. This Letter is the first application of the freezing trick to dynamical quantities.

In the U(2) spin CS model, Uglov has derived the exact formula of the dynamical spin-density correlation function

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with a finite number of particles [17]. We shall first extend his result to the case of $K \geq 3$, and then take the strong coupling limit. In doing so we have to make correspondence between physical quantities defined in the continuum and discrete spaces. Let us define the following operator in the continuum space:

$$X_q^{\gamma\delta} = \frac{1}{\sqrt{L}} \sum_{j=1}^{N} X_j^{\gamma\delta} e^{-iqx_j}, \tag{4}$$

where the momentum q takes values $2\pi n/L$ with n an arbitrary integer. We first derive the dynamical structure factor in the continuum model defined by

$$S^{(\gamma\delta)}(q,\omega;\beta) = \sum_{\alpha} |\langle\alpha|X_q^{\gamma\delta}|0\rangle|^2 \delta(\omega - E_\alpha + E_0), \tag{5}$$

where $\{|\alpha\rangle\}$ is the normalized complete basis of the system with eigenvalues $\{E_\alpha\}$, and $|0\rangle$ is the ground state. We assume that N is an integer multiple of K so that the ground state is a nondegenerate singlet.

In the strong coupling limit the coordinate x_j in Eq. (4) is written as $x_j = R_j + u_j$ where $R_j = j$ is a lattice point and u_j describes the lattice vibration. Except for the uniform motion of the lattice described by $u_j =$ const, we may regard u_j as a small quantity. In fact, the density response can be shown to be smaller than the spin response by $\mathcal{O}(\beta^{-1})$. Then the dynamical structure factor of the SU(K) HS model is given simply by the strong coupling limit of Eq. (5) provided that one restricts q in the range of the first Brillouin zone: $|q| \leq \pi$.

The dynamical structure factor (5) can be derived in a manner analogous to the case of $K = 2$ [17]. However, the following observations are necessary for generalization. First, each excited state relevant to Eq. (5) transforms as one of the weight vectors for the adjoint representation of SU(K). This observation allows us to find the selection rule for the SU(K) spin. In the SU(2) case, this selection rule is reduced to the simple fact that excited states relevant to Eq. (5) are spin-triplet states. Second, in order to derive the matrix element in Eq. (5) we find a convenient set of operators given by

$$J_a = \frac{1}{\sqrt{L}} \sum_q \left(\sum_{b=1}^{a} X_q^{b,K-a+b} + \sum_{b=1}^{K-a} X_q^{a+b,b} \right) \tag{6}$$

with $a = 1,\ldots,K-1$. More details of the calculation will be presented elsewhere.

In order to give the formula of $S^{(\gamma\delta)}(q,\omega)$, we fix some notations for partitions [17]. Let $\Lambda_N^{(K)}$ be the set of all partitions whose lengths are less than $N+1$ and the largest entry is less than $K+1$. Namely we have $\Lambda_N^{(K)} = \{\lambda = (\lambda_1,\lambda_2,\ldots,\lambda_N) \mid K \geq \lambda_1 \geq \lambda_2 \geq \ldots \geq \lambda_N \geq 0\}$. For a partition λ, we define subsets by $C_K(\lambda) = \{(i,j) \in \lambda \mid j - i \equiv 0 \bmod K\}$ and $H_K(\lambda) = \{(i,j) \in \lambda \mid \lambda_i + \lambda_j' - i - j + 1 \equiv 0 \bmod K\}$. Next we define the concept *type* of a partition. The *type* of a given partition λ contains sufficient information for determining the sets $C_K(\lambda)$ and $H_K(\lambda)$ explicitly. We introduce a reductive transformation τ on the set of all partitions as follows [22]: (i) If there exist K rows or K columns which have the same number of boxes in a partition, remove those rows or columns. (ii) Apply the reduction (i) repeatedly until the newly generated partition is no longer reducible.

We then determine a subset $\mathcal{A}_N^{(K)}$ of $\Lambda_N^{(K)}$ as the image of τ, i.e., $\mathcal{A}_N^{(K)} = \tau(\Lambda_N^{(K)})$. For any partition $\lambda \in \Lambda_N^{(K)}$, we say that ν is of the *type* λ if $\nu = \tau(\lambda)$. The total number of *types* increases from 3 in the case of SU(2) to 25 in SU(3), and to 252 in SU(4). For a box $s = (i,j) \in \lambda$, the numbers $l_\lambda(s) = \lambda_j' - i$ and $l_\lambda'(s) = i - 1$ are called the leg length and coleg length, respectively. For any subset $\nu \subset \lambda$, the order $|\nu|$ is defined as the number of boxes in ν.

Now we give the exact formula for the dynamical structure factor $S^{(\gamma\delta)}(q,\omega)$ for a finite size of the system with $\gamma \neq \delta$. We obtain

$$S^{(\gamma\delta)}(q,\omega) = \sum_{\lambda}{}' |F_\lambda^{(K)}|^2 \delta(\omega - E_\lambda), \tag{7}$$

where the primed summation is restricted so as to satisfy the momentum conservation $q = 2\pi|C_K(\lambda)|/N$ and the color selection rule for $\lambda \in \Lambda_N^{(K)}$. The latter is conveniently implemented by introducing a subset $\mathcal{A}_N^{(K;\gamma\delta)}$ of $\mathcal{A}_N^{(K)}$, and decomposes the summation over λ by each *type* $\nu = \tau(\lambda)$ such that $\nu \in \mathcal{A}_N^{(K;\gamma\delta)}$. For example, in the case of $K = 3$, we have $\mathcal{A}_N^{(3;21)} = \{(2,1,1),(3,3,1),(3,3,2,2)\}$ and $\mathcal{A}_N^{(3;13)} = \{(1),(2,2),(3,2,1,1)\}$.

In Eq. (7) the excitation energy is given by

$$E_\lambda = \frac{J}{4}\left(\frac{2\pi}{N}\right)^2 \left[(N-1)|C_K(\lambda)| - 2\sum_{s\in C_K(\lambda)} l_\lambda'(s)\right], \tag{8}$$

and the squared form factor by

$$|F_\lambda^{(K)}|^2 = \frac{1}{N} \frac{\prod_{s\in C_K(\lambda)/\{(1,1)\}} l_\lambda'(s)^2}{\prod_{s\in H_K(\lambda)} l_\lambda(s)[l_\lambda(s)+1]} \times \prod_{s\in C_K(\lambda)} \frac{N - l_\lambda'(s)}{N - l_\lambda'(s) - 1}. \tag{9}$$

Since we consider the case of zero external magnetic fields, the SU(K) symmetry demands that $S^{(\gamma\delta)}(q,\omega)$ is actually independent of (γ,δ) as long as $\gamma \neq \delta$. We can prove this fact using the expressions (7)–(9). In the particular case of $K = 2$, our formulas (7)–(9) generalize the known one [15] to the arbitrary size of the system. We have checked the validity of Eqs. (7)–(9) with $K = 2$ and 3 by comparing with the numerical result for $N \leq 24$ and $N \leq 15$, respectively. The numerical result is obtained via exact diagonalization and the recursion method. The agreement is excellent in both cases of $K = 2$ and 3. In Fig. 1, we present the result for $K = 3$ and $N = 15$.

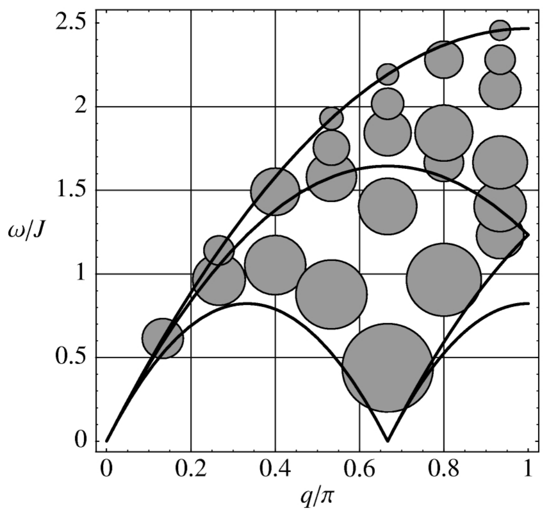

FIG. 1. Numerical result of the dynamical structure factor $S(q,\omega)$ in the cases of $K = 3$ and $N = 15$. The vertical and horizontal axes represent the rescaled energy and momentum, respectively. The intensity is proportional to the area of the circle. The solid lines are the dispersion lines of the elementary excitations in the thermodynamic limit. The analytic results are in excellent agreement with numerical ones, and are not distinguishable from the latter.

We now consider the quasiparticle interpretation of the color selection rule. For labeling the excited states relevant to $S^{(\gamma\delta)}(q,\omega)$, it is more convenient to use the conjugate partition $\lambda' = (\lambda'_1, \ldots, \lambda'_K) \in \Lambda_K^{(N)}$ instead of $\lambda \in \Lambda_N^{(K)}$. Each λ'_i has the information on the momentum and SU(K) spin of a quasiparticle. We call this quasiparticle a spinon following the SU(2) case. The spinon is considered to be an object possessing the SU(K) spin. Here the SU(K) spin means the $K - 1$ eigenvalues $(s_1, \ldots, s_{K-1})$ of a set of operators $(h^1, \ldots, h^{K-1})$ where h^γ is defined by $h^\gamma = \sum_{i=1}^N (X_i^{\gamma\gamma} - X_i^{\gamma+1\gamma+1})/2$ for $\gamma = 1, \ldots, K-1$. For the SU(2) case, this definition gives the z component of spin. The SU(K) spin of the spinon with λ'_i is specified by a certain condition on the pair (i, λ'_i). Since the condition for general K is rather complicated, we give an example in the case of $K = 3$. The SU(3) spin for λ'_i is assigned as follows: It is $(0, 1/2)$ if $(i, \lambda'_i) \equiv (0,0), (1,1), (2,2) \bmod 3$; $(1/2, -1/2)$ if $(i, \lambda'_i) \equiv (0,1), (1,2), (2,0) \bmod 3$; and $(-1/2, 0)$ if $(i, \lambda'_i) \equiv (0,2), (1,0), (2,1) \bmod 3$. It is important to note that a spinon transforms as a weight vector of the fundamental representation $\bar{K}$ of SU(K).

From the formulas (7)–(9), we can conclude that the relevant excited states for the SU(K) HS model consist of K spinons. Moreover the conditions on the *type* of excited states lead to an important consequence: K-spinon excited states have $K - 1$ different SU(K) spins. That is, the excited states contain $K - 1$ species of quasiparticles. For instance, in the case of $K = 3$, excited states relevant to $S^{(13)}(q,\omega)$ consist of three spinons with SU(3) spins $(0, 1/2)$, $(0, 1/2)$, and $(1/2, -1/2)$. For the SU(2) case, we recover the well-known fact that only two spinons with the same spin contribute to $S(q,\omega)$ [15].

The K-spinon excitation belongs to the tensor representation $\bar{K}^{\otimes K}$ of SU(K). This representation contains the adjoint representation as an irreducible component. From the condition stated above for the K-spinon excitation, we see that the K-spinon excitation transforms as one of the weight vectors for the adjoint representation. This is consistent with the condition for the SU(K) spins of the excited states which are relevant to $S^{(\gamma\delta)}(q,\omega)$.

Now we present the thermodynamic limit of the formulas (7)–(9). Performing a procedure similar to that in Refs. [23,24], for $K = 2$, 3, and 4, we obtain the final result as follows:

$$S(q,\omega) = A_K \sum_{1 \le a < b \le K} \prod_{i=1}^K \int_{-1}^1 dk_i \, |F_{ab}^{(K)}(k)|^2 \times \delta[q - \pi - p(k)] \delta[\omega - \epsilon(k)], \quad (10)$$

where $\epsilon(k) = [\pi v_s/(2K)] \sum_{i=1}^K (1 - k_i^2)$ with $v_s = J\pi/2$, and $p(k) = (\pi/K) \sum_{i=1}^K k_i$. In the above formula, A_K is a normalization constant given by

$$A_K = \frac{2^K \pi}{K^3(K-1)} \prod_{j=1}^K \frac{\Gamma[(K-1)/K]}{\Gamma(j/K)^2}, \quad (11)$$

and the form factor is given by

$$F_{ab}^{(K)}(k) = \frac{|k_a - k_b|^{g_K} \prod_{1 \le i < j \le K, (i,j) \ne (a,b)} |k_i - k_j|^{g'_K}}{\prod_{i=1}^K (1 - k_i^2)^{(1-g_K)/2}} \quad (12)$$

with $g_K = (K-1)/K$ and $g'_K = -1/K$. Since the formula (10) for $S^{(\gamma\delta)}(q,\omega)$ does not depend on the pair (γ, δ) with $\gamma \ne \delta$, we have omitted the superscript. Unfortunately our exact result is inconsistent with a conjecture proposed several years ago [25].

For $K = 2$, the formula reproduces the result of Haldane-Zirnbauer [15] which was obtained by a completely different method. Notice that the second product in the numerator of Eq. (12) is absent in the SU(2) case.

We can derive static structure factor $S(q)$ by integrating over ω in Eq. (10). In the low energy limit we recover the results of Ref. [9] which are obtained by conformal field theory. By analyzing the form factor for $q \sim 2k_F$ with $k_F = \pi/K$, we can show that it has the asymptotic form

$$S(q) \sim a_1 |q - 2k_F|^{\alpha_1 - 1} \quad (13)$$

with the exponent $\alpha_1 = 2 - 2/K$ and a nonuniversal constant a_1. In the real space the spin (or color) correlation decays as $b_1 \cos(2k_F x)|x|^{-\alpha_1}$ with a certain coefficient b_1. A similar analysis shows that there are also weaker singularities around $q = 2lk_F$ for $l = 2, \ldots, K-1$ with exponents $\alpha_l = 2l(1 - l/K)$. These $K - 1$ singularities correspond to $K - 1$ gapless bosonic modes [7,9].

The spinon interpretation of the formula (10) goes as follows: As in the case of finite systems, the excited states for

$S(q,\omega)$ in the thermodynamic limit consist of K spinons with $K-1$ different SU(K) spins. This fact means, as in the low energy limit, the dynamics of the SU(K) HS model can be described by the $K-1$ species of quasiparticles. This simple structure reflects the Yangian symmetry of the SU(K) HS model [26]. In the form factor (12), the factor $|k_a - k_b|^{g_K}$ represents the statistical interactions of spinons with the same SU(K) spin, while the factor $|k_i - k_j|^{g'_K}$ represents those of spinons with different SU(K) spins. We refer to Ref. [27] for a more detailed explanation of statistical interactions. It will be interesting to consider the relation between our results and the exclusion statistics in conformal field theory discussed in Refs. [28,29].

The support of $S(q,\omega)$ represents the region in the momentum-frequency plane where $S(q,\omega)$ takes the nonzero value. We see that the support of $S(q,\omega)$ as determined from the formula (10) is compact, i.e., there is no intensity outside of the finite area. In the SU(3) case, for example, the support is determined as

$$\omega \le [v_s/(2\pi)]q(2\pi - q) \equiv \epsilon^{(\mathrm{U})}(q) \quad \text{for } 0 \le q \le 2\pi\,,$$

$$\omega \ge [3v_s/(2\pi)]q(2\pi/3 - q) \quad \text{for } 0 \le q \le 2\pi/3\,,$$

$$\omega \ge [3v_s/(2\pi)](q - 2\pi/3)(4\pi/3 - q)$$
$$\text{for } 2\pi/3 \le q \le 4\pi/3\,,$$

$$\omega \ge [3v_s/(2\pi)](q - 4\pi/3)(2\pi - q)$$
$$\text{for } 4\pi/3 \le q \le 2\pi\,.$$

For a general value of K, there are K lower boundaries as given by

$$\epsilon_j^{(\mathrm{L})}(q) \equiv [Kv_s/(2\pi)](q - 2\pi j/K)[2\pi(j+1)/K - q]\,,$$

for $2\pi j/K \le q \le 2\pi(j+1)/K$ with $j = 0,1,\ldots,K-1$.

The behavior of $S(q,\omega)$ near the boundaries of the support is derived for general K as follows: We can show that there is a stepwise discontinuity at the upper boundary $\omega = \epsilon^{(\mathrm{U})}(q)$. On the other hand, there are divergent singularities at the lower boundaries $\omega = \epsilon_0^{(\mathrm{L})}(q)$ and $\omega = \epsilon_{K-1}^{(\mathrm{L})}(q)$. Here $S(q,\omega)$ diverges by the power law with the exponent $-1/K$. At the other lower boundaries $\omega = \epsilon_j^{(\mathrm{L})}(q)$ with $j \ne 0, K-1$, $S(q,\omega)$ has threshold singularities but no divergence. As in the SU(2) case [15,30], we expect that the divergences at two of the lower boundaries occur also in the SU(K) Heisenberg model with the nearest-neighbor exchange.

In conclusion, we have derived the exact formulas (7) and (10) for $S(q,\omega)$ of the SU(K) HS model for the arbitrary size of the system at zero temperature. Our exact result of $S(q,\omega)$ for $K \le 4$ is likely to be valid for larger K as well. We have also clarified the quasiparticle picture of the spin dynamics. The relevant excited states consist of K spinons with $K-1$ different SU(K) spins.

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