Contents

TWO-DIMENSIONAL SPANNING WEBS AS (1,2) LOGARITHMIC MINIMAL MODEL

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An example of exactly solvable logarithmic models with conformal boundary conditions has recently been given by Pearce and Rasmussen [1]. They considered critical dense polymers with certain types of defects on a strip and reproduced the conformal weights in the first column of the extended Kac table.

The model we study is a generalization of the spanning tree model on a finite square lattice wrapped on a cylinder. We consider spanning webs in which noncontractible cycles are allowed. In the case of closed boundaries at both edges of the cylinder, the tree branches must be attached to the noncontractible cycles, while an open boundary allows for trees rooted at the sites of this boundary as well. The model is similar in many aspects to the model of critical dense polymers studied in [1]. The entries of the Kac table for the latter model are labeled by the number of defect lines under fixed boundary conditions at both sides of the strip. We show that the cycles in the spanning webs play the role of pairs of defect lines in the model of dense polymers, see Fig. 1.

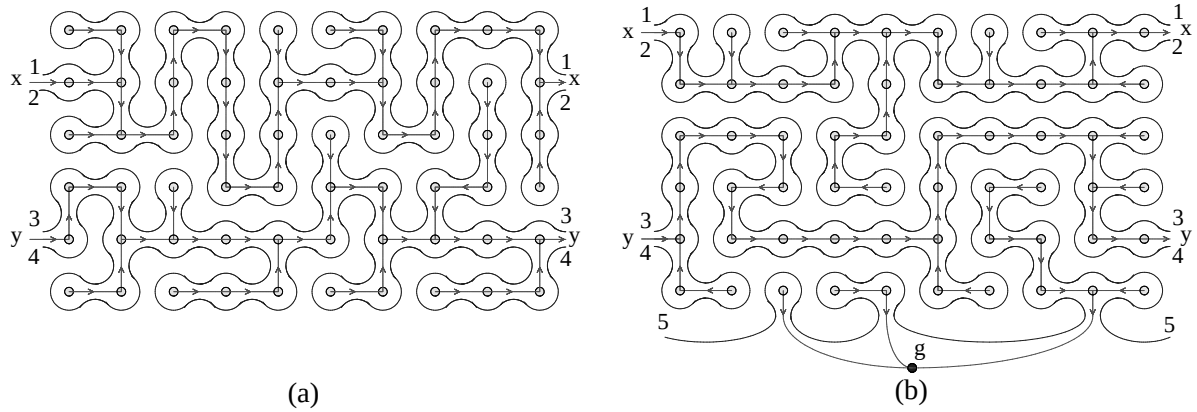

Рис. 1: Correspondence between spanning webs and dense polymers. (a) The cycle $x - x$ is surrounded by two polymers, $1 - 1$ and $2 - 2$, and the cycle $y - y$ by another two polymers, $3 - 3$ and $4 - 4$ which are considered as defect lines in the classification of [1, 2]. (b) Besides the four defect polymer lines, there is a fifth line $5 - 5$ separating the boundary trees from the rest of the lattice.

The spanning web model differs from that of the dense polymers in several essential aspects. The cylinder geometry admits parametrization of the web configurations by the number of noncontractible cycles which are well defined 'quantum numbers'. The finite-size partition function of our model is calculated by using an extension of the Kirchhoff theorem and evaluation of standard determinant expressions for the free fermion model. The most important result is the perfect coincidence of the universal part of the partition function for all the different combinations of closed and open boundaries with the finitized characters of symplectic fermions under periodic or antiperiodic boundary conditions.

The exact analytical expression for the partition function of the spanning web model

on a cylinder with perimeter M and length N is given by:

$$Z_{NM}^{(\mu,\nu)} = \prod_{p=0}^{N-1} [Q_N^{(\mu,\nu)}(p)]^{2M} \left\{ 1 + (\omega + \omega^{-1})[Q_N^{(\mu,\nu)}(p)]^{-2M} + [Q_N^{(\mu,\nu)}(p)]^{-4M} \right\}. \tag{1}$$

Here,

$$\begin{aligned} &Q \equiv Q_N^{(\mu,\nu)}(p) = \sqrt{1 + \sin^2 \phi_N^{(\mu,\nu)}(p)} + \sin \phi_N^{(\mu,\nu)}(p), \\ &\phi_N^{(0,0)}(p) = \frac{\pi p}{2N}, \quad \phi_N^{(0,1)}(p) = \phi_N^{(1,0)}(p) = \frac{\pi(2p+1)}{2(2N+1)}, \quad \phi_N^{(1,1)}(p) = \frac{\pi(p+1)}{2(N+1)}, \end{aligned} \tag{2}$$

the parameter ω is the weight of a cycle encircling the cylinder in the positive direction, and ω^{-1} is the corresponding weight for the opposite direction. The superscripts $\mu, \nu = 0, 1$, fix the type of boundary conditions at the top and bottom edges of the cylinder: $\mu = 0$ stays for closed and $\mu = 1$ for open boundary. The asymptotic form of the free energy, as $N \to \infty$ and $M \to \infty$, so that N/M remains fixed, separates into terms including the bulk and surface contributions,

$$F_{MN}^{(\mu,\nu)} = \frac{4G}{\pi} MN + M \left[\frac{2G}{\pi}(\mu + \nu) - \ln(1 + \sqrt{2}) \right], \tag{3}$$

and the *universal part*:

$$\ln \bar{Z}_N^{(\mu,\nu)}(q, \omega) = \ln \left(q^{\frac{1}{12} - \frac{1}{8}(\mu-\nu)^2} \prod_{j=0}^{N-1} \left[1 + (\omega + \omega^{-1}) q^{j + \frac{1}{2}(\mu+\nu)} + q^{2j+\mu+\nu} \right] \right). \tag{4}$$

Here G is Catalan's constant and $q = \exp(-\pi M/N)$. Expression (4) follows from the approximation $Q_N^{(\mu,\nu)}(p) \simeq 1 + \phi_N^{(\mu,\nu)}(p)$.

In the absence of cycles, the finite-size effects in the spanning tree model, particularly the Casimir effect, show that it corresponds to a logarithmic minimal model with central charge $c = -2$ [3]. The presence of cycles changes the Casimir effect in accordance with the conformal weights which appear in the Kac table [2]. Moreover, the partition function (4) can be interpreted in terms of symplectic fermions [5]. The symplectic fermions are fermionic fields $\theta^{\pm}(z)$ with operator product expansion

$$\theta^+(z)\theta^-(w) \sim \log(z - w). \tag{5}$$

These fields admit periodic and antiperiodic boundary conditions under which they decompose with integer $\theta_n^{\pm}$, $n \in \mathbb{Z}$, and half-integer $\theta_n^{\pm}$, $n \in \mathbb{Z} + \frac{1}{2}$, modes, respectively. These modes satisfy the anticommutation relations

$$[\theta_n^+, \theta_m^-]_+ = n\delta_{n+m,0}. \tag{6}$$

Let $\mathcal{A}(2)$ denote this infinite dimensional Clifford algebra. The algebra $\mathcal{A}(2)$ has two irreducible modules X_1 and X_2 (see details in [4, 6]). In addition, the algebra $\mathcal{A}(2)$ has two projective modules P_1 and $P_2 = X_2$. The module P_1 contains 4 irreducible subquotients isomorphic to X_1 [6]. The character of X_1 is

$$\chi^{(1,1)}(q, \omega) = \frac{q^{\frac{1}{12}}}{\prod_{n=1}^{\infty}(1 - q^n)} \sum_{r \in \mathbb{N}} \sum_{j=0}^{r} \omega^{r-2j} q^{\frac{r(r-1)}{2}} (1 - q^r). \tag{7}$$

The character of X_2 is

$$\chi^{(0,1)}(q,\omega) = \frac{q^{-\frac{1}{24}}}{\prod_{n=1}^{\infty}(1-q^n)} \sum_{r\in\mathbb{N}} \sum_{j=0}^{r} \omega^{r-2j} q^{\frac{(r-1)^2}{2}} (1-q^{2r}). \tag{8}$$

The character of P_1 is

$$\chi^{(0,0)}(q,\omega) = (2+\omega+\omega^{-1})\chi^{(1,1)}(q,\omega). \tag{9}$$

We fix the subalgebra $\mathcal{A}(2)[N]$ of $\mathcal{A}(2)$ for $N \in \mathbb{N}$

$$\mathcal{A}(2)[N] = \begin{cases} \{\theta^{\pm}_{-n},\ n \geq N\} & \text{periodic b. c.} \\ \{\theta^{\pm}_{-n-\frac{1}{2}},\ n \geq N\} & \text{antiperiodic b. c.} \end{cases} \tag{10}$$

and consider the characters $\chi^{(1,1)}[N](q,\omega)$, $\chi^{(0,1)}[N](q,\omega)$ and $\chi^{(0,0)}[N](q,\omega)$ of coinvariants with respect to $\mathcal{A}(2)[N]$ in the modules X_1, X_2 and P_1, respectively. The series expansion in ω of the product in the universal part of the partition function (4) has the same form as Eqs. (7) - (9). Moreover, the following equality for the finitized characters holds:

$$\chi^{(\mu,\nu)}[N](q,\omega) = \bar{Z}_N^{(\mu,\nu)}(q,\omega). \tag{11}$$

Thus, universal part of the partition function $\bar{Z}_N^{(\mu,\nu)}(q,\omega)$ for the spanning web model under different combinations of closed and open boundary conditions coincides with the finitized characters of the symplectic fermions. This allows us to interpret the simplectic fermion model as a conformal field theory of spanning webs on a cylinder.

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SELF-CONSISTENT RENORMALIZATION THEORY OF SPIN FLUCTUATIONS IN THE METALLIC SPINEL LiV_2O_4

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1. In the family of transition metal oxides the spinel compound LiV_2O_4 is a rare metallic system showing heavy fermion behavior. In particular, an anomalously large specific heat coefficient $\gamma = C/T$ and strongly enhanced magnetic susceptibility χ_s were detected in the low temperature limit, $T < 30$K. Despite continuous activity in the last years, there is currently no consensus on the mechanism for formation of heavy fermion quasiparticles in LiV_2O_4, and the issue is still under debate. Recently, we have suggested [1, 2, 3] our own route for explaining unusual properties of this material.

LiV_2O_4 has the cubic spinel structure with the magnetic vanadium ions (in the mixed valence state $V^{3.5+}$) occupying the pyrochlore lattice sites. At low temperatures, $T < 30$ K, the spin system of LiV_2O_4 exhibits pronounced short-range antiferromagnetic (AFM) correlations, but no long-range magnetic ordering was detected at any measured temperatures. The geometrical frustration of the pyrochlore lattice is likely to be a crucial aspect of the problem. The frustration may suppress at any T a long-range ordering of strongly correlated itinerant electrons, but instead, the system is placed near a magnetic instability. The emergence of largely degenerate low-lying spin excitations in the ground state of LiV_2O_4 is expected to be responsible for low-T properties of this material, including its heavy fermion behaviour. This appealing picture was developed [1] and further examined [2, 3] by comparing theoretical results with experimental data obtained by different techniques, like the inelastic neutron scattering (INS) and the nuclear magnetic resonance (NMR), probing low frequency spin fluctuations.

2. Recalling that LiV_2O_4 is found to be a paramagnet down to very low temperatures, we assumed that the system remains spin disordered at all temperatures and calculated the dynamic spin susceptibility $\chi(\mathbf{q},\omega)$ in the ground state at $T = 0$. The calculations were performed with the use of the electronic band structure obtained for LiV_2O_4 in the local density approximation (LDA) and effects of strong electron correlations were treated in the random phase approximation (RPA). We found [1] that in the low-ω limit the calculated $\chi_{RPA}(\mathbf{q},\omega \to 0)$ was picked at the wave vectors of length $|\mathbf{q}| \sim |\mathbf{Q}_c| \sim 0.6$ Å^{-1}, i.e., at the same position where INS intensity exhibits the maximum of a broad distribution. This remarkable coincidence indicates that the location in $\mathbf{q}$-space and the multiplicity of "critical"wave vectors $\mathbf{Q}_c$ are produced by specific properties of the electronic band structure and the topology of the many sheet Fermi surface in LiV_2O_4.

3. To extend the theory for $T > 0$, one step beyond the RPA was made and a self-consistent renormalization (SCR) theory of spin fluctuations in LiV_2O_4 was developed [2]. The SCR theory offers a phenomenological description for "critical"spin fluctuations in nearly antiferromagnetic itinerant electron systems by taking into account effects of mode-mode coupling between spin fluctuations at $|\mathbf{q}| \sim |\mathbf{Q}_c|$. In this theory, the temperature dependent inverse static susceptibility at $|\mathbf{q}| = |\mathbf{Q}_c|$ defined as $y_Q(T) = [2T_A\chi(Q_c,\omega = 0;T)]^{-1}$ is of principal importance.

First, we derived [2] that the inverse susceptibility $y_Q(t)$, where $t = T/T_0$ was the reduced temperature, obeyed the following integral equation:

$$y_Q(t) = y_Q(0) + g_Q \int_0^{z_c} dz \frac{\phi\left(\left[y_Q(t) + z^2\right]/t\right) - \phi\left(\left[y_Q(t) + z^2 + bx_c\right]/t\right)}{bx_c/t}, \quad (1)$$

with

$$\phi(u) = \ln\Gamma(u) - \left(u - \frac{1}{2}\right)\ln u + u - \frac{1}{2}\ln 2\pi, \quad (2)$$

where $\Gamma(u)$ was the gamma function.

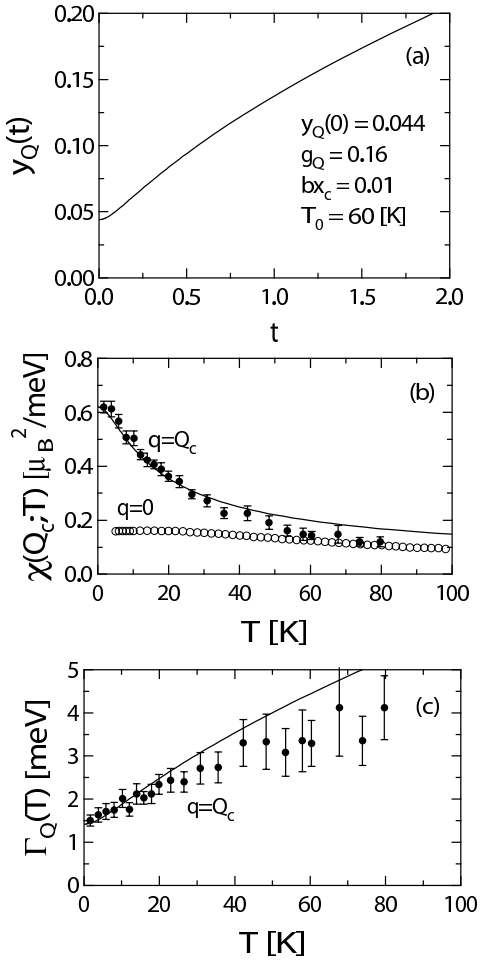

Рис. 1: (a) The solution of Eq.(1) for the inverse static spin susceptibility at $q = Q_c$ as a function of the reduced temperature $t = T/T_0$; the full set of fit parameters is given in the text; (b) Static spin susceptibilities at $q = Q_c$ (solid circles) and $q = 0$ (open circles) as functions of temperature observed in INS and magnetic measurements on LiV_2O_4. The solid line is a fit to $\chi(Q_c;T)$ using the same solution as in (a); (c) Spin relaxation rate $\Gamma_{Q_c}(T)$ at $q = Q_c$ as a function of temperature observed in INS experiment. The solid line is a fit using again the solution of (a). The INS data are taken from S.-H. Lee.*et al* [4].

In the present form, the SCR theory is parametrized with five parameters which are now denoted as $y_Q(0)$, T_A, T_0, g_Q and bx_c. The parameters T_A and T_0 characterize, as $T \to 0$, the momentum and frequency spread of "critical"spin fluctuations, g_Q is the effective mode-mode coupling constant and bx_c is a measure of the anisotropy of the spin fluctuation dispersion in **q**-space. At the final stage, we put the theory on a quantitative ground by adjusting the parameter values when comparing the calculated model results with the available experimental data, including INS and magnetic measurements for the spin susceptibility in LiV_2O_4.

The best overall fit to experimental data is achieved with the following parameters: $T_0 \simeq 60K, T_A \simeq 220K, g_Q = 0.16, bx_c = 0.01$ and $y_Q(0) \simeq 0.044$. As shown, for instance, in Fig.1, the experimental INS data are well represented by our theory. As expected, a better agreement between the theory and experiment is reached in the low temperature

region, $T < 40$K. Actually, for higher temperatures the AFM fluctuations at $|\mathbf{q}| \simeq Q_c$ are suppressed and no more distinguished from those at other wave vectors; the system enters gradually a spin localized regime compatible with the Curie-Weiss behavior of $\chi(\mathbf{q}=0,\omega=0)$ observed in LiV_2O_4 for $T > 60$ K.

4. The NMR measurements of the spin-lattice relaxation rate $1/T_1$ as a function of temperature and external pressure is another experimental method for studying spin dynamics in LiV_2O_4. Based on the SCR theory, our estimate [3] for the contribution of the "critical"spin fluctuation to the spin-lattice relaxation rate is as follows:

$$\left(\frac{1}{T_1T}\right)_{q\sim Q_c} = \frac{3\gamma_n^2\hbar|A_{Q_c}|^2}{\pi k_B T_0 T_A}\left(\frac{Q_c}{q_B}\right)^2\frac{1}{bx_c}\left\{\frac{1}{\sqrt{y_Q(T)}}\tan^{-1}\frac{z_c}{\sqrt{y_Q(T)}} - \frac{1}{\sqrt{y_Q(T)+bx_c}}\tan^{-1}\frac{z_c}{\sqrt{y_Q(T)+bx_c}}\right\}, \quad (3)$$

where A_{Q_c} is the hyperfine coupling constant and the reduced inverse susceptibility $y_Q(t)$ obeys the integral equation (1). We consider $y_Q(0)$ to be the only pressure-dependent parameter, assuming that $y_Q(0) \to 0$ as the system approaches a quantum critical point with increasing pressure.

The experimental data [5] for $(1/T_1T)$ measured for different applied pressures together with theoretical curves [3] fitting these data are depicted in Figure 2. A good agreement between the theoretical results and the available experimental data is found at least down to 1 K.

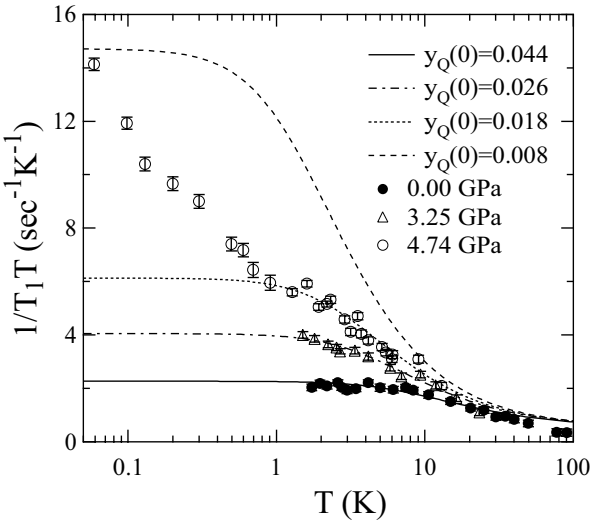

Рис. 2: Temperature dependence of $1/T_1T$ obtained from ^{7}Li-NMR measurements on powder samples of LiV_2O_4 under different applied pressure; the data are taken from K. Fujiwara *et al*, [5]. Different curves together with the corresponding values of the fit parameter $y_Q(0)$ represent the calculations based on the SCR theory.

The SCR theory fails, however, in giving a quantitative description of the low-T upturn of $1/T_1T$ detected below 0.6 K under the highest applied pressure of 4.74 GPa. The discrepancy can be clearly explained by taking into account that a certain amount of crystal defects and/or magnetic impurities that are unavoidably present in the measured powder samples of LiV_2O_4 may contribute to the nuclear spin relaxation as well. Actually, when approaching a magnetic instability and softening of largely degenerate low lying spin fluctuations, the system becomes very susceptible to weak perturbations including, for instance, magnetic defects.

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APPROXIMATE EXPRESSION FOR THE DYNAMIC STRUCTURE FACTOR IN THE LIEB-LINIGER MODEL

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Cigar-shaped traps with cold alkali atoms have recently been used to obtain a quasi-1D quantum degenerate Bose gas, where atomic motion in the transverse dimensions is confined to zero-point quantum oscillations, in weak and strong interaction regimes [1]. Theoretically, we can describe this system as a one-dimensional gas where interactions of bosonic atoms can be described well by effective δ-function interactions [2]. Thus, the Lieb-Liniger model is applicable. Being exactly solvable in the uniform case, the model, however, does not admit complete analytic solutions for the correlation functions. Up to now, this has been a challenging and still unsolved problem in 1D physics. Here, we propose an approximate formula for the DSF of the Lieb-Liniger gas that is consistent with the known results in accessible limits and power laws.

Dynamical density-density correlations, which can be measured by the two-photon Bragg scattering [3], are described by the *dynamic structure factor* (DSF)

$$S(q,\omega) = L \int \frac{\delta t \delta x}{2\pi\hbar}\, e^{i(\omega t - qx)} \langle 0|\delta\hat{\rho}(x,t)\delta\hat{\rho}(0,0)|0\rangle. \tag{1}$$

Here, we introduce the density fluctuations $\delta\hat{\rho}(x,t) \equiv \hat{\rho}(x,t) - n$ and the equilibrium density of particles $n = N/L$. We consider the case of zero temperature, where $\langle 0|\ldots|0\rangle$ denotes the ground-state average.

The Lieb-Liniger model represents a uniform 1D system of spinless bosons of mass m, interacting with pairwise point interactions $V(x) = g_{\mathrm{B}}\delta(x)$; the interaction strength g_{B} is assumed to be positive. The strength of interactions can be measured in terms of the dimensionless Lieb-Liniger parameter $\gamma \equiv mg_{\mathrm{B}}/(\hbar^2 n)$. Within the Lieb-Liniger model, the DSF has the following well-established properties.

i) Luttinger liquid theory predicts a power-law behaviour of the DSF at low energies in the vicinity of the momenta $k = 0, 2\pi n, 4\pi n \ldots$ and yields universal values of the exponents [4]. In particular, one can show [4] that in the vicinity of "umklapp"point ($k = 2\pi n$, $\omega = 0$), $S(k,\omega) \sim (\omega^2 - \omega_-^2)^{K-1}$, where $K \equiv \hbar\pi n/(mc)$ and c is the sound velocity. Furthermore, within the Luttinger-liquid theory, the dispersion is linear in the vicinity of the umklapp point: $\omega_-(k) \simeq c|k - 2\pi n|$. This relation leads to the exponent $2(K-1)$ precisely at the umklapp point and $K-1$ outside of it.

ii) By using in a nontrivial manner the Bose-Fermi mapping in 1D, the authors developed the time-dependent Hartree-Fock scheme [5] in the *strong-coupling* regime with a small parameter $1/\gamma$. The scheme guarantees the validity of the DSF expansion [5]

$$S(k,\omega)\frac{\varepsilon_{\mathrm{F}}}{N} = \frac{k_{\mathrm{F}}}{4k}\left(1 + \frac{8}{\gamma}\right) + \frac{1}{2\gamma}\ln\frac{\omega^2 - \omega_-^2}{\omega_+^2 - \omega^2} + O\left(\frac{1}{\gamma^2}\right), \tag{2}$$

for $\omega_- \leqslant \omega \leqslant \omega_+$, and zero otherwise. Here $\omega_\pm(k)$ are the limiting dispersions that bound quasiparticle-quasihole excitations in the linear response theory. In the strong-coupling regime, they are given by $\omega_\pm(k) = \hbar|2k_{\mathrm{F}}k \pm k^2|(1 - 4/\gamma)/(2m) + O(\gamma^{-2})$. By definition,

$k_F \equiv \pi n$ and $\varepsilon_F \equiv \hbar^2 k_F^2/(2m)$ are the Fermi wave vector and energy of a noninteracting Fermi gas, respectively.

iii) As was shown by Imambekov and Glazman [6], in the Lieb-Liniger model the DSF exhibits a power-law behaviour near the borders $\omega_\pm(k)$

$$S(k,\omega) \sim |\omega - \omega_\pm(k)|^{\mp\mu_\pm(k)}. \tag{3}$$

The positive exponents $\mu_\pm$ are related to the quasiparticle scattering phase and can be easily evaluated by solving a system of a few integral equations in the thermodynamic limit. The exact relation $\mu_-(2\pi n - 0) = 2\sqrt{K}(\sqrt{K}-1)$ can be obtained [6], which obviously differs from the Luttinger liquid exponent $2(K-1)$ for $k \neq 2\pi n$. However, the Imambekov and Glazman result is correct in the immediate vicinity of $\omega_\pm$ provided that the finite curvature of $\omega_-(k)$ is taken into account. Thus, the difference in the exponents can be treated [6] as an artifact of the linear spectrum approximation in the Luttinger liquid theory. Note, however, that the thin "strip"in the ω-k plane where the exponents are different vanishes at the point $k = 2\pi n$; hence, the Luttinger exponent $2(K-1)$ should be exact here.

iv) The DSF can be calculated numerically with the algebraic Bethe ansatz [7].

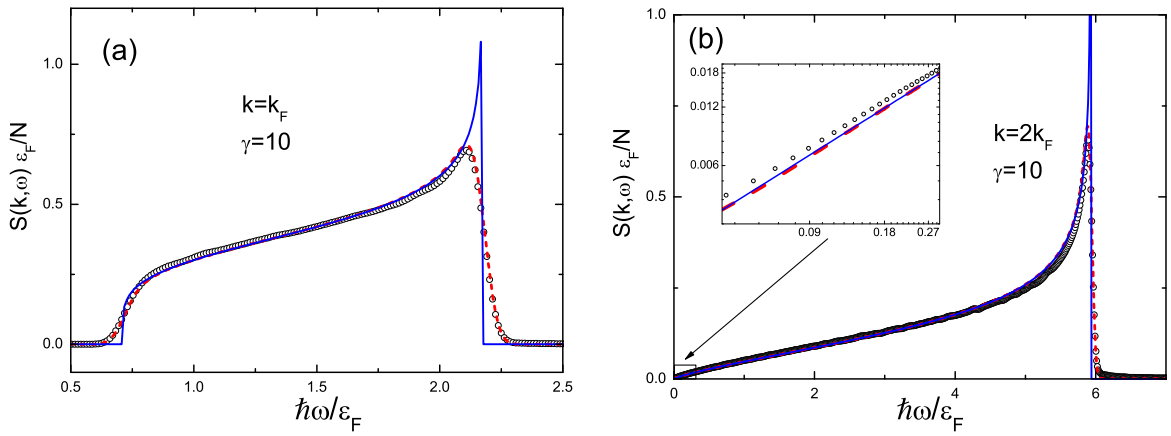

Рис. 1: DSF in the thermodynamic limit. The proposed approximation (4) (line) is compared to numerical data [7] (open dots). The dashed (red) line represents the data of Eq. (4) convoluted in frequency with a Gaussian of width $0.07\varepsilon_F/\hbar$ in order to simulate smearing that was used in Ref. [7] while generating the numerical results. The numerical data [7] suggest that contributions from multi-particle excitations for $\omega > \omega_+$ (sharp line in parts **a** and **b**) are very small. Such contributions are not accounted for by the formula (4). Insert: DSF at the umklapp point in logarithmic scale. The graph shows that the DSF behaves as predicted by the Luttinger liquid theory with the exponent $2(K-1)$, where $K = 1.402\ldots$ at $\gamma = 10$.

We suggest a phenomenological expression [8] which is consistent with all of the above-mentioned results. It reads

$$S(k,\omega) = C\frac{(\omega^\alpha - \omega_-^\alpha)^{\mu_-}}{(\omega_+^\alpha - \omega^\alpha)^{\mu_+}} \tag{4}$$

for $\omega_-(k) \leqslant \omega \leqslant \omega_+(k)$, and zero otherwise. Here C is a normalization constant, $\mu_+(k)$ and $\mu_-(k)$ are the exponents of Eq. (3), and $\alpha \equiv 1 + 1/\sqrt{K}$. The normalization constant depends on the momentum but not the frequency and can be determined from the f-sum

rule $\int_{-\infty}^{+\infty} \delta\omega\, \omega S(q,\omega) = Nq^2/(2m)$. We assume that in Eq. (4) the value of the exponent $\mu_-(2\pi n)$ coincides with its limiting value $\mu_-(2\pi n - 0) = 2\sqrt{K}(\sqrt{K} - 1)$.

Now it can be easily seen from Eq. (4) that the exponent equals $2(K-1)$ at the umklapp point and $\mu_-(k)$ outside of it. Thus, the suggested formula is consistent with both the Luttinger liquid behaviour at the umklapp point and the Imambekov and Glazman power-law behaviour in the vicinity of it, as it should be. In the strong-coupling regime, Eq. (4) correctly yields the first order expansion (2). Besides, comparison with numerical data by Caux and Calabrese [7] (Fig. 1) shows that the suggested formula nicely works in the regimes of both the weak and the strong coupling.

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ENTANGLEMENT IN COUPLED BOSON SYSTEMS

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There is currently an enormous effort underway to understand the rich dynamics of interacting Bose systems. Among various problems are the properties of Bose-Einstein condensate (BEC) in external fields [1], the Bose coherent effects of excitons and polaritons in semiconductor microcavities [2], and the production of scattered radiation due to interaction of the incident laser beam with vibrational modes of a medium (the Raman effect). The generic feature of interacting Bose systems is the formation of Bose-field collective states with nonclassical (squeezed, sub-Poissonian, etc.) statistical and fluctuation properties. Another prominent feature is an entanglement produced by the interaction between different system constituents, which is one of the most subtle and intriguing phenomena in nature (see for a recent review [3]). Nowadays, there is explosive activity in the study of the entanglement due to its potential usefulness in quantum teleportation, quantum cryptography, and, in general, in quantum information theory.

The dynamical interplay between quantum entanglement and nonclassical properties of various Bose systems can be traced within a model of two-coupled harmonic oscillators. This model has been applied in nuclear physics and for a rotating BEC. In condensed matter physics this sort of a Hamiltonian is used: i) to study the interaction between an atom and a radiative field; ii) as a starting point for analysis of electronic properties of two-dimensional quantum dots in a perpendicular magnetic field. The dynamics of two-component Bose condensate trapped in a double-well potential can also be mapped on the time-evolution of two coupled-harmonic oscillators in the low excitation regime. In the simplest case, the model describes two levels of the condensed atoms which are coupled owing to the classical field of radiation. The dynamic interplay between the strength of the interaction, entanglement and squeezing dynamics of the system was analyzed in [4] in the general case assuming that at the initial stage of the time evolution the fields were found in the superposition of coherent and chaotic states. It can correspond to the Bose condensate of “cold"atoms being described by a coherent state at a nonzero temperature which spoils the plain coherence and brings the decoherence effects. The Gaussian character of the system state allows one to calculate analytically the measure of entanglement between the fields in the form of the logarithmic negativity via the symplectic spectrum of the partial transpose of the covariance matrix. It is shown that the interaction with equal coupling constants produces the entangled system state (see Fig. 1a) when the degree of entanglement displays an oscillatory character with the distinctive alteration of the entanglement maxima and minima that can be considered to be revivals and collapses. Thermal fluctuations are expected to attenuate the entanglement. In particular, when both the fields are initially in chaotic states, the degree of their entanglement in the regime of equal coupling constants becomes very small (see Fig. 1b). Moreover, there is a critical temperature above which the entanglement disappears and the system becomes separable. Thus, the considered two-mode boson model demonstrates that the entanglement can be controlled in time by a appropriate choice of the interaction at different temperature regimes. The obtained results show that this plain system exhibited a rich dynamics from the view point of the quantum information theory and its possible physical applications.

Another process for generation of quantum correlations is the Raman scattering. The Stokes and anti-Stokes fields in the Raman scattering, which indirectly interact via a

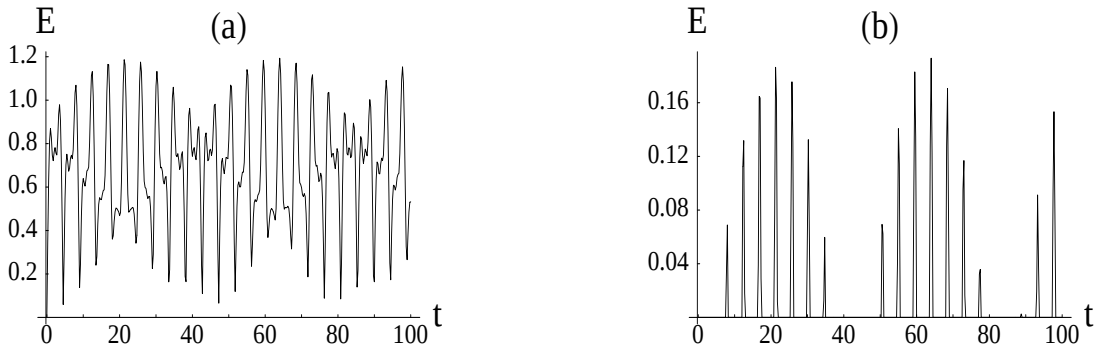

Рис. 1: Entanglement between the modes being initially in (a) coherent states and (b) chaotic states for the interaction with equal coupling constants.

bath of phonons, can be modelled by two nondegenerate boson fields coupled through a reservoir. The appearance of quantum entanglement between these fields with regard to initial states of the fields and the phonon bath state was analyzed in [5]. It was shown that there occurred a phenomenon of entanglement between the Stokes and anti-Stokes fields due to their correlation through the phonon bath in the case of stimulating scattering, when these fields are initially found in superpositions of coherent and chaotic states, and the phonon bath is in a chaotic state. The beginning of entanglement for initial coherent states of the fields is revealed to arise from the very outset of the interaction. The strength of entanglement monotonously increases in time approaching its asymptotic value. In the case of initial chaotic states, the entanglement can arise only after a certain interaction time. The noise of the phonon reservoir appeared to work for destruction of entanglement tending to reduce its strength. Moreover, if the phonon noise is large enough, then the entanglement might not happen between the fields at all. Therefore, the phonon bath in the model comes to be a source of both quantum correlations and their decoherence.

Summarizing, the peculiar pattern of entanglement between the boson fields found out in systems under consideration is expected to have promising applications in quantum communication and quantum information processing.

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THE GEOMETRY-INDUCED EFFECTS CAUSED BY POSITIVE AND NEGATIVE DISCLINATIONS IN CARBON NANOSTRUCTURES

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The two main features of carbon nanostructures are their effective two-dimensional behaviour and topological defects (namely-disclinations) included in it. Our model [1] includes both disclinations and curvature of the surface describing the electronic properties of carbon nanostructures within the field-theoretic approach.

The electronic states of capped semi-infinite nanotubes are studied in our work [2] within the phenomenological gauge field-theory model. A single manifold for the description of both the nanotube and the cap region (considered as nearly a half of either Ih or I fullerene) is suggested. For the cap region, the gauge fields which take into account six disclinations are taken from the previously studied model for spherical fullerenes [3]. The wavefunctions and the density of states (DoS) are numerically calculated for both metallic and semiconducting nanotubes. The smoothing of van Hove singularities is found (see Fig.1) and proven analytically. The effect of smoothing itself is a very

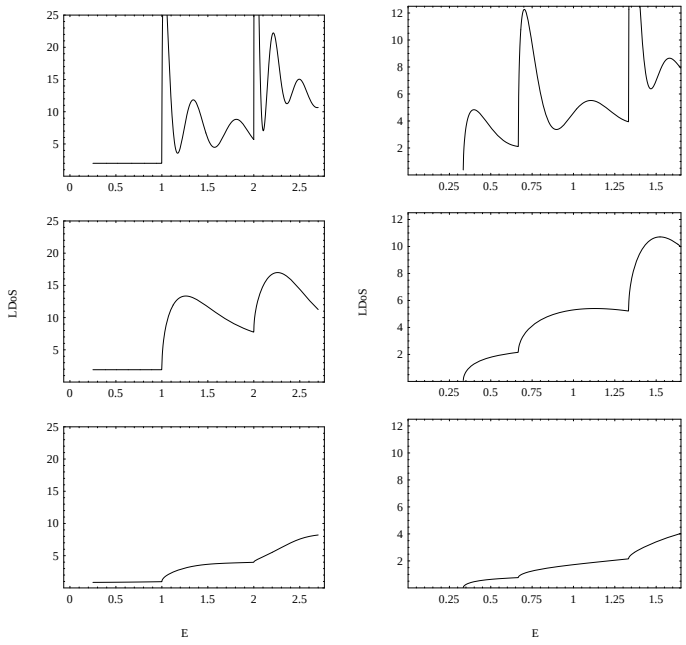

Рис. 1: Local density of states (average units) on the cap (bottom), near the cap and far from the cap (top), for metallic (left) and semiconducting (right) nanotubes. The energy E is measured in units of $\hbar V_F/R_t$, where V_F- Fermi speed, R_t- nanotube radius.

interesting phenomenon, appearing usually as a result of electron-electron or electron-phonon interaction. Within our model, however, this effect appears as a result of purely geometrical factors (i.e., the change of the nanotube radius).

In [4], we studied the electronic structure of graphene in the presence of either sevenfolds or eightfolds by using a gauge field-theoretic model. The graphene sheet with

topological defects is considered as a "negative cone surface"with an infinite Gaussian curvature at the center (see Fig.2); the surface is supposed to be free to bend but impossible to stretch. The density of electronic states is calculated for a single seven-

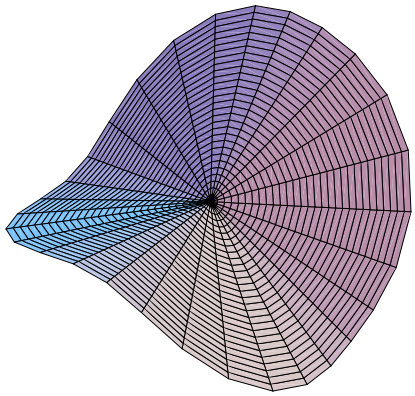

Рис. 2: A surface of the graphene sheet with a single sevenfold inserted in it.

and eightfold as well as for a pair of sevenfolds with different morphology. Because of the non-Abelian character of the gauge field, the gauge field for two closely positioned sevenfolds is not equal to the field of one eightfold, as it is in the usual elasticity theory. The translational factor M should be taken into account. The density of states at the Fermi level is found to be zero in all cases except two sevenfolds with translational factor $M \neq 0$.

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ELLIPTIC HYPERGEOMETRIC FUNCTIONS AND SUPERCONFORMAL INDICES

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One of the important recent achievements of mathematical physics consists in the discovery of elliptic hypergeometric functions – a principally new class of special functions of hypergeometric type (see [1] for a survey of the corresponding results and relevant literature). These functions have found applications in the theory of Yang-Baxter equation, integrable discrete time chains, elliptic Calogero-Sutherland type models and so on. Although connection with the classical root systems has been explicitly traced in the structure of many elliptic hypergeometric functions, their group theoretical content remained obscure by large.

In recent papers, Römelsberger [2] has described topological indices for four dimensional supersymmetric conformal field theories. Following [2], Dolan and Osborn [3] connected superconformal indices of a number of $\mathcal{N}=1$ supersymmetric field theories with certain elliptic hypergeometric integrals. The corresponding Seiberg dual theories [4] appeared to have the same indices due to nontrivial identities for these integrals [1].

Seiberg duality is a fundamental concept of modern quantum field theory [4, 5]. We have performed a systematic comparison of superconformal indices of known dual supersymmetric theories and elliptic hypergeometric integrals. In [6], we considered $\mathcal{N}=1$ theory with $SP(2N)$ gauge group and fixed rank flavor groups $SU(8)$ or $SU(4)\times SU(4)\times U(1)_B$. For $N=1$ we obtained indices for multiple dual theories considered by Csáki et al [5]. For $N>1$ we came to new multiple duality theories.

We start from the electric theory with the overall symmetry group

$$SP(2N)\times SU(8)\times U(1)\times U(1)_R.$$

The field content for the electric theory is given in the table

	$SP(2N)$	$SU(8)$	$U(1)$	$U(1)_R$
Q	f	f	$N-1$	$2R_Q=\frac{1}{2}$
V	adj	1	0	1
X	T_A	1	-4	0

Following the Römelsberger algorithm, we find that the superconformal index in electric theory is given by the integral

$$I_E = \frac{(p;p)_\infty^N (q;q)_\infty^N}{2^N N!}\Gamma((pq)^s;p,q)^{N-1}\int_{\mathbb{T}^N}\prod_{1\le i<j\le N}\frac{\Gamma((pq)^s z_i^{\pm 1}z_j^{\pm 1};p,q)}{\Gamma(z_i^{\pm 1}z_j^{\pm 1};p,q)} \times\prod_{j=1}^N\frac{\prod_{1\le i\le 8}\Gamma((pq)^{r_Q}y_i z_j^{\pm 1};p,q)}{\Gamma(z_j^{\pm 2};p,q)}\prod_{j=1}^N\frac{dz_j}{2\pi i z_j}, \tag{1}$$

where

$$\Gamma(z;p,q)=\prod_{j,k=0}^\infty\frac{1-z^{-1}p^{j+1}q^{k+1}}{1-zp^jq^k}$$

is the elliptic gamma function, s is an arbitrary parameter, $r_Q = R_Q - (N-1)s/4$, and the balancing condition has the form $\prod_{j=1}^{8} y_j = 1$.

We skip consideration of the $N = 1$ case, which appears to be related to the E_7-group symmetry transformation for an elliptic analogue of the Gauss hypergeometric function discovered by the first author [1]. As shown in [6], for $N > 1$ one can construct several dual magnetic theories for the initial electric theory applying the Rains symmetry transformation discussed in [1] for integral (1). The first magnetic theory has the overall symmetry group

$$SP(2N) \times SU(4) \times SU(4) \times U(1)_B \times U(1) \times U(1)_R$$

with the following spectral content of the fields:

	$SP(2N)$	$SU(4)$	$SU(4)$	$U(1)_B$	$U(1)$	$U(1)_R$
q	f	f	1	-1	$N-1$	$\frac{1}{2}$
$\widetilde{q}$	f	1	f	1	$N-1$	$\frac{1}{2}$
$\widetilde{V}$	adj	1	1	0	0	1
Y	T_A	1	1	0	-4	0
$M_j, j = 0, \ldots, N-1$	1	T_A	1	2	$2N-2-4j$	1
$\widetilde{M}_j, j = 0, \ldots, N-1$	1	1	T_A	-2	$2N-2-4j$	1

The index in the magnetic theory is then

$$\begin{aligned} I_M^{(1)} \;=\; & \prod_{l=0}^{N-1} \prod_{1\le i<j\le 4} \Gamma((pq)^{r_{M_l}} y_i y_j; p, q) \prod_{5\le i<j\le 8} \Gamma((pq)^{r_{\widetilde{M}_l}} y_i y_j; p, q) \\ & \times \frac{(p;p)_\infty^N (q;q)_\infty^N}{2^N N!} \Gamma((pq)^s; p, q)^{N-1} \int_{\mathbb{T}^N} \prod_{1\le i<j\le N} \frac{\Gamma((pq)^s z_i^{\pm 1} z_j^{\pm 1}; p, q)}{\Gamma(z_i^{\pm 1} z_j^{\pm 1}; p, q)} \\ & \times \prod_{j=1}^{N} \frac{\prod_{1\le i\le 4} \Gamma((pq)^{r_q} y_i z_j^{\pm 1}; p, q) \prod_{5\le i\le 8} \Gamma((pq)^{\widetilde{r}_q} y_i z_j^{\pm 1}; p, q)}{\Gamma(z_j^{\pm 2}; p, q)} \prod_{j=1}^{N} \frac{dz_j}{2\pi i z_j}. \end{aligned} \tag{2}$$

Using the Rains transformation one more time we come to the second magnetic dual theory which has the same overall group as in the previous case but different representation content

	$SP(2N)$	$SU(4)$	$SU(4)$	$U(1)_B$	$U(1)$	$U(1)_R$
q	f	$\overline{f}$	1	1	$N-1$	$2R_q = \frac{1}{2}$
$\widetilde{q}$	f	1	$\overline{f}$	-1	$N-1$	$2\widetilde{R}_q = \frac{1}{2}$
$\widetilde{V}$	adj	1	1	0	0	1
Y	T_A	1	1	0	-4	0
$M_j, j = 0, \ldots, N-1$	1	f	f	0	$2N-2-4j$	1

The index for the described magnetic theory equals

$$I_M^{(2)} \;=\; \frac{(p;p)_\infty^N (q;q)_\infty^N}{2^N N!} \Gamma((pq)^s; p, q)^{N-1} \prod_{l=0}^{N-1} \prod_{\substack{1\le i\le 4 \\ 5\le j\le 8}} \Gamma((pq)^{r_{M_l}} y_i y_j; p, q)$$

$$\times \int_{\mathbb{T}^N} \prod_{1\le i<j\le N} \frac{\Gamma((pq)^s z_i^{\pm 1} z_j^{\pm 1};p,q)}{\Gamma(z_i^{\pm 1} z_j^{\pm 1};p,q)} \tag{3}$$
$$\times \prod_{j=1}^{N} \frac{\prod_{1\le i\le 4}\Gamma((pq)^{r_q} y_i^{-1} z_j^{\pm 1};p,q)\prod_{5\le i\le 8}\Gamma((pq)^{\widetilde{r}_q} y_i^{-1} z_j^{\pm 1};p,q)}{\Gamma(z_j^{\pm 2};p,q)} \prod_{j=1}^{N}\frac{dz_j}{2\pi i z_j}.$$

Finally, the third magnetic theory has the overall symmetry group

$$SP(2N)\times SU(8)\times U(1)\times U(1)_R.$$

The field content is

	$SP(2N)$	$SU(8)$	$U(1)$	$U(1)_R$
q	f	$\overline{f}$	$N-1$	$\frac{1}{2}$
$\widetilde{V}$	adj	1	0	1
Y	T_A	1	-4	0
$M_j,\ j=0,\dots,N-1$	1	T_A	$2N-2-4j$	1

The magnetic index is now

$$I_M^{(3)} = \frac{(p;p)_\infty^N (q;q)_\infty^N}{2^N N!}\Gamma((pq)^s;p,q)^{N-1}\prod_{l=0}^{N-1}\prod_{1\le i<j\le 8}\Gamma((pq)^{r_{M_l}} y_i y_j;p,q) \tag{4}$$
$$\times \int_{\mathbb{T}^N}\prod_{1\le i<j\le N}\frac{\Gamma((pq)^s z_i^{\pm 1} z_j^{\pm 1};p,q)}{\Gamma(z_i^{\pm 1} z_j^{\pm 1};p,q)}\prod_{j=1}^{N}\frac{\prod_{1\le i\le 8}\Gamma((pq)^{\widetilde{r}} y_i^{-1} z_j^{\pm 1};p,q)}{\Gamma(z_j^{\pm 2};p,q)}\prod_{j=1}^{N}\frac{dz_j}{2\pi i z_j}.$$

So we have obtained three dual magnetic theories for a fixed $\mathcal{N}=1$ electric supersymmetric theory with the $SP(2N)$ gauge group and a fixed number of flavors.

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