Y-system and Quasi-Classical Strings

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ABSTRACT: Recently Kazakov, Vieira and the author conjectured the Y-system set of equations describing the planar spectrum of AdS/CFT. In this paper we solve the Ysystem equations in the strong coupling scaling limit. We show that the quasi-classical spectrum of string moving inside $AdS_3 \times S^1$ matches precisely with the prediction of the Y-system. Thus the Y-system, unlike the asymptotic Bethe ansatz, describes correctly the spectrum of one-loop string energies including all exponential finite size corrections. This gives a very non-trivial further support in favor of the conjecture.

KEYWORDS: AdS/CFT, Integrability.

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1. Introduction

Since the discovery of AdS/CFT correspondence [1, 2] there has been a significant progress in solving the maximally super-symmetric Yang-Mills theory in four dimensions, mainly due to the integrability found on both sides of the duality [3, 4, 5, 6, 7]. The problem of finding the anomalous dimensions of all local operators in planar limit was thus brought very close to its complete solution. Recently the anomalous dimension of the simplest operator was computed numerically for a wide range of 't Hooft coupling λ [8].

The approach used there is based on the Y-system for AdS/CFT conjectured in [9]. The main ingredients leading to the conjecture are the asymptotic Bethe ansatz equations (ABA) [10, 11], describing the spectrum of the anomalous dimensions of local operators with infinitely many constituent fields, the concept of the "mirror" double Wick rotated theory invented in [12] and then explored in details in [13], Lüscher formula for the finite



Figure 1: T-shaped "fat hook". It defines the interactions between Y-functions in the AdS/CFT Y-system equations.

size corrections [14] and our experience with relativistic theories, where similar equations are constantly appearing in the thermodynamic Bethe ansatz (TBA) approach [15, 16].

This conjecture passes some nontrivial tests – in [9] the 4-loop perturbative result [17] was reproduced. More recently a comparison was also made at 5-loops in [18]. In [19, 20] it was also shown to be consistent with the TBA approach and finally its numerical solution [8] indicates that the strong coupling asymptotic agrees with the string prediction [2] for the simplest Konishi operator¹. Apparently more tests of this conjecture are indispensable. In this paper we compare the results of the quasi-classical string quantization with the prediction of Y-system thus providing the first analytical test of the conjecture to all orders in finite size wrapping corrections. As we will see below this check involves numerous remarkable identities and miraculous simplifications and probes most of the Y-system in great detail. Therefore, the match we will observe leaves little room for doubt about the validity of the Y-system, at least in the strong coupling limit.

The Y-system of [9] is an infinite system of simple functional equations²

$$Y_{as}(z+\frac{i}{4g})Y_{as}(z-\frac{i}{4g}) = \frac{(1+Y_{a,s+1}(z))(1+Y_{a,s-1}(z))}{(1+1/Y_{a+1,s}(z))(1+1/Y_{a-1,s}(z))},$$
(1.1)

The indices a, s belong to a T-shaped lattice (see Fig.1). One should replace Y_{as} with indices outside the lattice by 0 or ∞ so that they disappear from the equations. Note that the Y-system equations only involve the pair of indices a, s and its immediate neighbors only, $a \pm 1, s$ and $a, s \pm 1$; in this sense these equations are described as "local". In accordance with Fig.1 we also use the following notations³

$$\{Y_{\bullet_a}, Y_{\triangle_a}, Y_{\bigcirc_s}, Y_{\otimes}, Y_{\oplus}\} = \{Y_{a,0}, Y_{a,1}, Y_{1,s}, Y_{1,1}, Y_{2,2}\}.$$
(1.2)

²In this paper we use a rescaled rapidities $z = \frac{u}{2g}$, where $g = \frac{\sqrt{\lambda}}{4\pi}$, more convenient for our purposes.

¹The sub-sub-leading coefficient does not agree with the quasi-classical string results of [21]. However, the Konishi state is one of the lowest states and the quasi-classical quantization may not be applicable for this case.

 $^{^{3}}$ In the original derivation of the Y-system equations the Y-functions are associated with densities of several bound states of the mirror theory [22]. The symbols used as subscripts in this new notation were introduced in [19] to indicate which type of bound-state they originate from.

To have a unique solution this system should be supplemented with an additional "non-local" equation $[19]^4$

$$\log Y_{\otimes}Y_{\oplus} = \sum_{a=1}^{\infty} \left(\mathcal{R}^{(0a)} - \mathcal{B}^{(0a)} \right) * \log(1 + Y_{\bullet_a}) + \log \frac{R^{(+)}B^{(-)}}{R^{(-)}B^{(+)}}, \qquad (1.3)$$

and particular boundary conditions at infinite a and s. Here, \mathcal{R} and \mathcal{B} are some kernels defined in Appendix A and * stands for convolution. This additional data encodes information about particular operator/state.

When the solution is found the energy can be computed from

$$E = \sum_{j=1}^{M} \epsilon_1^{\rm ph}(z_j) + \sum_{a=1}^{\infty} \int_{-\infty}^{\infty} \frac{dz}{2\pi i} \frac{\partial \epsilon_a^{\rm mir}(z)}{\partial z} \log(1 + Y_{\bullet_a}^{\rm mir}) , \qquad (1.4)$$

where $\epsilon_a(z)$ is a single magnon energy⁵ It is a simple multi-valued function and ph/mir indicate its branches (see Appendix A for notations). Similarly $Y_{\bullet_a}^{\text{ph}}$ and $Y_{\bullet_a}^{\text{mir}}$ correspond to the different branches. Finally z_i are exact Bethe roots satisfying the exact Bethe equation

$$Y_{\bullet_1}^{\rm ph}(z_i) = -1$$
 . (1.5)

In this paper we will focus on the $\mathfrak{sl}(2)$ subsector of the correspondence. From the gauge theory point of view this subsector is a closed subsector, with operators being composed from scalar fields \mathcal{Z} and covariant derivatives \mathcal{D} ,

$$\operatorname{Tr}\left(\mathcal{D}^{M}\mathcal{Z}^{L}\right) + \dots$$
 (1.6)

The dots stand for all possible permutations of the derivatives with \mathcal{Z} fields. In the Ysystem language, M is the number of roots z_j whereas L enters through large a asymptotic of Y_{\bullet_a} . From the string theory side, this subsector describes strings moving in $AdS_3 \times S^1$ contained in the larger $AdS_5 \times S^5$ space-time. At infinite coupling the string motion is classical while the first $1/\sqrt{\lambda}$ corrections are the quasi-classical one-loop effects which we study in this paper. For this subsector the solution is $s \to -s$ symmetric $Y_{a,s} = Y_{a,-s}$, which makes our consideration more transparent.

In the limit when the number of operators L goes to infinity, $Y_{\bullet_a}^{\min}$ are exponentially suppressed and the integral term in (1.4) becomes irrelevant. The remaining term is the sum of the individual energies of magnons. In this limit the solution of Y-system is known for arbitrary coupling $g = \frac{\sqrt{\lambda}}{4\pi}$ and arbitrary state [9]. The finite size effects enter in two different ways: firstly, the second term in (1.4) becomes important and, secondly, the Bethe roots z_i move away from their asymptotic values giving rise to the modification of the first term.

⁴We strongly believe that Hirota equation associated with this Y-system should automatically imply (1.3). From this point of view Hirota equation seems to be more fundamental. See also discussion in [9].

⁵Magnons are the fundamental spin chain excitations around the BPS vacuum tr Z^L . They are dual to the string worldsheet excitations. Their dispersion relation and the scattering matrix describing magnon scattering are known at any coupling and are properties of the infinite volume theory.

In this paper we solve the Y-system equations (1.1) in the strong coupling limit. We construct explicitly the Y-functions for large number of Bethe roots $M \sim g \gg 1$. Then we compare the results with the quasi-classically quantized string using methods developed in [23, 24, 25, 26, 27, 29, 30] and show their perfect match. We show that the Y-system resolves the know disagreement of asymptotic Bethe ansatz with the semi-classically quantized strings [31] due to finite-size effects.

2. Y-system in the scaling limit

The scaling limit is the strong coupling limit $g \to \infty$ where the string can be described classically. In this limit, the number of Bethe roots M and the operator length L go to infinity as g. We also assume that the Bethe roots $z_i \sim 1$ and $|z_i| > 1$. This limit of Bethe equations was introduced in [32, 33]. The Bethe roots form continuous cuts in the complex plane z. They become the branch-cuts on the classical algebraic curve which we describe in Sec.3.1.

In the scaling limit we can neglect i/(4g) shifts in the arguments in l.h.s. of (1.1) which with $1/g^2$ precision becomes a set of algebraic equations. This is the key simplification which will allow us to completely solve the Y-system at strong coupling. In the notations (1.2) we have three infinite series of equations

$$Y_{O_s}^2 = (1 + Y_{O_{s+1}})(1 + Y_{O_{s-1}}) , \quad s = 3, 4, \dots ,$$
(2.1)

$$Y_{\Delta_a}^2 = \frac{(1+Y_{\bullet_a})}{(1+1/Y_{\Delta_{a+1}})(1+1/Y_{\Delta_{a-1}})} , \ a = 3, 4, \dots ,$$
 (2.2)

$$Y_{\bullet_a}^2 = \frac{(1+Y_{\triangle_a})^2}{(1+1/Y_{\bullet_{a+1}})(1+1/Y_{\bullet_{a-1}})} , \quad a = 2, 3, \dots ,$$
 (2.3)

plus four equations

$$Y_{\Delta_2}^2 = \frac{(1+Y_{\oplus})(1+Y_{\bullet_2})}{(1+1/Y_{\Delta_3})(1+1/Y_{\otimes})}, \qquad (2.4)$$

$$Y_{O_2}^2 = \frac{(1+Y_{O_3})(1+Y_{\otimes})}{(1+1/Y_{\oplus})}, \qquad (2.5)$$

$$Y_{\otimes}^{2} = \frac{(1+Y_{\bigcirc_{2}})(1+Y_{\bullet_{1}})}{(1+1/Y_{\triangle_{2}})}, \qquad (2.6)$$

$$Y_{\bullet_1}^2 = \frac{(1+Y_{\otimes})^2}{(1+1/Y_{\bullet_2})} \,. \tag{2.7}$$

From the asymptotic solution of [9], which we re-consider in the next section, one expects that Y-functions have several branch cuts going from $z = \pm 1 + \frac{in}{4g}$, for some integers n, to infinity. The approximation we use above is only accurate far enough from the branch-cuts. Close to the branch-cuts even a small shift in the argument could cause a large jump of Y's. This means we can safely use (2.1-2.7) above and below the real axis and on the interval [-1, 1] of the real axis, but not on the whole real axis. In this section we will mainly focus on the spectral parameters with -1 < Re z < 1 for which (2.1-2.7) are valid. Equation (1.3) also simplifies in the scaling limit

$$\frac{R^{(+)}B^{(-)}}{R^{(-)}B^{(+)}} = \prod_{j=1}^{M} \frac{x(z) - x_j^-}{x(z) - x_j^+} \frac{1/x(z) - x_j^+}{1/x(z) - x_j^-} \simeq \frac{1}{f(z)\bar{f}(z)} , \qquad (2.8)$$

where

$$x(z) = z + i\sqrt{1-z^2}$$
, $x_j = x^{\text{ph}}(z_j) = z_j + \sqrt{z_j - 1}\sqrt{z_j + 1}$ (2.9)

and $x_j^{\pm} = x^{\text{ph}}(z_j \pm \frac{i}{4g})$ for the standard definition of $\sqrt{\ldots}$ with branch cut along negative part of the real axis. The functions f(z) and $\bar{f}(z)$ are defined in terms of the resolvent

$$G(x) = \frac{1}{g} \sum_{j}^{M} \frac{1}{x - x_j} \frac{x_j^2}{x_j^2 - 1} , \qquad (2.10)$$

in the following way

$$f(z) = \exp\left(-iG(x(z))\right) , \quad \overline{f}(z) = \exp\left(+iG(1/x(z))\right). \quad (2.11)$$

The kernels $\mathcal{R}^{(0a)}$ and $\mathcal{B}^{(0a)}$, defined in Appendix A simplify dramatically in the strong coupling scaling limit. For example $\mathcal{R}^{(0a)}(z,w) - \mathcal{B}^{(0a)}(z,w) \simeq \delta(z-w)$ so that equation (1.3) becomes

$$F \equiv Y_{\otimes}Y_{\oplus} = \frac{1}{f\bar{f}} \prod_{a=1}^{\infty} \log(1 + Y_{\bullet_a}) .$$
(2.12)

One should keep in mind that the possible distribution of the Bethe roots is constrained by

$$\sum_{j=1}^{M} \frac{1}{i} \log \frac{x_j^+}{x_j^-} \simeq \sum_{j=1}^{M} \frac{x_j}{g(x_j^2 - 1)} = 2\pi m + \mathcal{O}(1/g) \quad , \quad m \in \mathbb{Z} \; , \tag{2.13}$$

which reflects the cyclicity symmetry of the single trace operators.

We see that in the scaling limit we have to solve an infinite set of algebraic equations (2.1-2.7,2.12). In the next section we construct the asymptotic solution of these equations for $Y_{\bullet_a} \ll 1$ and expand it in the scaling limit.

2.1 Infinite length solution at strong coupling

In this section we will study the Y-system in the asymptotic limit $Y_{\bullet_a} \ll 1$ and then expand it at strong coupling. We will see that in this scaling limit the asymptotic solutions of [9] can be recast in the following simple form:

$$Y_{\bigcirc_s}(z) = (s-A)^2 - 1 \quad , \quad Y_{\triangle_a} = \frac{(T-1)^2 S T^{a-1}}{(ST^{a+1} - 1)(ST^{a-1} - 1)} \quad , \tag{2.14}$$

where

$$A = \frac{1}{\bar{f} - 1} + \frac{f}{f - 1} , \quad S = \frac{\bar{f}(f - 1)^2}{f(\bar{f} - 1)^2} , \quad T = \frac{f}{\bar{f}} .$$
 (2.15)

with (2.11). The middle node Y-functions are equally simple:

$$Y_{\bullet_a} = \Delta^a \left(\frac{f(\bar{f}-1)^2 \bar{f}^a - (f-1)^2 \bar{f} f^a}{f \bar{f}(\bar{f}-f)} \right)^2 , \qquad (2.16)$$

where

$$\Delta = \exp\left(-\frac{\frac{L}{2g} + 2\pi mz}{\sqrt{1 - z^2}}\right) , \qquad (2.17)$$

governs their exponential suppression for large L/g. We now derive these equations starting from the asymptotic solution constructed in [9] in terms of the eigenvalues of SU(2|2)transfer matrices for representations with rectangular Young tableaux T_{as} . We have

$$1 + 1/Y_{\bigcirc_s}(z) = \frac{T_{1,s}(z + \frac{i}{4g})T_{1,s}(z - \frac{i}{4g})}{T_{1,s-1}(z)T_{0,s+1}(z)} , \quad 1 + Y_{\triangle_a}(z) = \frac{T_{a,1}(z + \frac{i}{4g})T_{a,1}(z - \frac{i}{4g})}{T_{a-1,1}(z)T_{a+1,1}(z)} . \quad (2.18)$$

The middle node functions $Y_{\bullet_a}(z)$ are suppressed but can be also expressed in terms of the transfer matrices using

$$Y_{\bullet_{a}}(z) = (T_{a,1})^{2} \left(\frac{x(z-\frac{ia}{4g})}{x(z+\frac{ia}{4g})}\right)^{L} \prod_{b=-\frac{a-1}{2}}^{\frac{a-1}{2}} \prod_{j=1}^{M} \frac{B^{(+)}(z+i\frac{2b+1}{4g})R^{(-)}(z+i\frac{2b-1}{4g})}{B^{(-)}(z+i\frac{2b-1}{4g})R^{(+)}(z+i\frac{2b+1}{4g})} \sigma^{2}(z+\frac{ib}{2g},x_{j}) ,$$
(2.19)

where $B^{(\pm)}$, $R^{(\pm)}$ are defined in the Appendix A and $\sigma(z, x_j)$ for Im z > 0 coincides with the Beisert-Eden-Staudacher dressing phase [10] and is understood as an analytical continuation otherwise. The SU(2|2) transfer matrices for symmetric $(T_{1,s})$ and antisymmetric $(T_{a,1})$ representations can be found from the expansion of the generating functional [34, 35] (see [36] for some mathematical details). For the $\mathfrak{sl}(2)$ subsector it reads

$$\mathcal{W} = \left[1 - \frac{B^{(+)}(z + \frac{i}{4g})R^{(+)}(z - \frac{i}{4g})}{B^{(-)}(z + \frac{i}{4g})R^{(-)}(z - \frac{i}{4g})}\hat{D}\right] \left[1 - (-1)^m \frac{R^{(+)}(z - \frac{i}{4g})}{R^{(-)}(z - \frac{i}{4g})}\hat{D}\right]^{-2} \left[1 - \hat{D}\right] , \quad (2.20)$$

where $\hat{D} = e^{-i\frac{\partial z}{2g}}$ is a shift operator. One should expand (2.20) in powers of \hat{D} and commute them to the right

$$\mathcal{W} = \sum_{s=0}^{\infty} T_{1s} \left(z - i \frac{s-1}{4g} \right) \hat{D}^s , \ \mathcal{W}^{-1} = \sum_{a=0}^{\infty} (-1)^a T_{a1} \left(z - i \frac{a-1}{4g} \right) \hat{D}^a .$$
(2.21)

We want to construct Y_{as} using this construction in the scaling limit. At strong coupling the shift operator \hat{D} serves just as a formal expansion parameter D, since the shifts it creates are suppressed. In the notations of the previous section \mathcal{W} becomes

$$\mathcal{W} = \frac{\left(1 - D\bar{f}/f\right)(1 - D)}{\left(1 - D/f\right)^2} \,. \tag{2.22}$$

Now it is very easy to find the general expression for T_{1s} and T_{a1}

$$T_{1s} = \frac{(s-1)f\bar{f} - s(f+\bar{f}) + (s+1)}{f^s} , \ T_{a1} = (-1)^a \frac{f(\bar{f}-1)^2\bar{f}^a - \bar{f}(f-1)^2f^a}{f\bar{f}(\bar{f}-f)f^a} , \ (2.23)$$

which implies the above mentioned expressions (2.14) for the Y-functions. In the expression for the middle node Y's (2.19) it is enough to use the leading strong coupling expression for the dressing phase – the Arutyunov–Frolov–Staudacher (AFS) phase [6]

$$\sigma_a^2(z, x_j) \simeq \left(\frac{1 - 1/(x^-(z)x_j^+)}{1 - 1/(x^+(z)x_j^-)}\right)^2 \left(\frac{x^-(z)x_j^- - 1}{x^+(z)x_j^- - 1}\frac{x^+(z)x_j^+ - 1}{x^-(z)x_j^+ - 1}\right)^{4ig(z_j - z)}, \quad (2.24)$$

where $x^{\pm} = x(z \pm \frac{i}{4g})$. Expanding this expression at strong coupling for -1 < z < 1 we obtain $Y_{\bullet_a} \simeq (f^2 \Delta)^a T_{a1}^2$, or (2.16).

We end this section with an important comment which will be later used in Sec.2.5. We notice that for |z| > 1 the second factor in (2.19) is exponentially small even for finite L/g (notice that x(z) denotes the mirror branch $x(z) = z + i\sqrt{1-z^2}$)

$$\left(\frac{x(z-\frac{ia}{4g})}{x(z+\frac{ia}{4g})}\right)^L \simeq \frac{1}{x(z)^{aL}} , \qquad (2.25)$$

and thus the asymptotic solution is accurate for |z| > 1 even if for -1 < z < 1 it is significantly modified by the finite size effects. Let us stress that the exponential suppression we are discussing is much stronger than the usual finite size exponential suppression at strong coupling. The latter is suppressed for large L/g whereas (2.25) is suppressed for large Leven if L/g is finite and small.

In the next section we will analyze the large a and s limit of these asymptotic Y's and argue that the same asymptotics should be used even when the fine size effects are strong.

2.2 Boundary conditions

In this section we propose the boundary conditions which should be used to make the solution of the Y-system unique. For that we study the asymptotic large L solution considered in the previous section at large a or s and argue that the exact solution should have exactly the same behavior.

From (2.14) we see that Y_{\triangle_a} oscillates with *a* because $T(z) = f(z)/\bar{f}(z)$ is a pure phase. To have a well defined large *a* limit we shift the argument by -i0 then |T(z-i0)| > 1 and we get

$$\lim_{a \to \infty} \frac{\log Y_{\triangle_a}(z-i0)}{a} = \log \frac{\bar{f}}{f} \,. \tag{2.26}$$

Similarly

$$\lim_{a \to \infty} \frac{\log Y_{\bullet_a}(z - i0)}{a} = \log \left(\Delta f^2\right) . \tag{2.27}$$

Whereas Y_{\triangle_a} and Y_{\bullet_a} decrease exponentially with a, Y_{\bigcirc_s} behave as s^2 . The general solution of (2.1) with polynomial asymptotics is

$$Y_{O_s} = (s - A(z))^2 - 1 \tag{2.28}$$

for some A(z). When Δ is small and Y_{\bullet_a} are suppressed A(z) is given by its asymptotic value (2.15), otherwise it is some unknown function. In the next section we find its exact expression as a function of Δ .

Conditions (2.26), (2.27) and (2.28) could be seen to be consistent with the TBA equations for excited states proposed in [19] (see also [20]) Y_{Δ_a} should satisfy⁶

$$\log Y_{\triangle_a}(z-i0) = \mathcal{M}_{am} * \log(1+Y_{\bullet_m}) - K_{a-1,m-1} * \log(1+Y_{\triangle_m}) - K_{a-1} * \log(1+Y_{\otimes}) + \log \frac{R_a^{(+)}B_{a-2}^{(+)}}{R_a^{(-)}B_{a-2}^{(-)}} + \log \frac{R_a^{(+)}B_{a-2}^{(+)}}{R_a^{(-)}B_{a-2}^{(+)}} + \log \frac{R_a^{(+)}B_{a-2}^{(+)}}{R_a^{(+)}B_{a-2}^{(+)}} + \log \frac{R_a^{(+)}B_{a-2}^{(+)}}{R_a^{(+)}B_{a-2}^{(+)}}} + \log \frac{R_a^{(+)}B_{a-2}^{(+)}}{R_a^{(+)}B_{a-2}^{(+)}}} + \log \frac{R_a^{(+)}B_{a-2}^{(+)}}{R_a^{(+)}} + \log \frac{R_a^{(+)}B_{a-2}^{(+)$$

where \mathcal{M} and K are some kernels defined in Appendix A. At strong coupling the last term gives

$$\lim_{a \to \infty} \frac{1}{a} \log \frac{R_a^{(+)} B_{a-2}^{(+)}}{R_a^{(-)} B_{a-2}^{(-)}} \simeq \log \frac{\bar{f}}{f} , \qquad (2.29)$$

assuming that the other terms are not growing with a linearly this leads precisely to (2.26). Similarly, (2.27) and (2.28) could be justified from the TBA equations of [19].

2.3 Y-system in *T*-hook

In this section we solve (2.1-2.7) together with (2.12). One can achieve a considerable simplification of this problem by transforming (2.1-2.7) into the Hirota equation. For that we rewrite Y_{as} in terms of T_{as}

$$Y_{a,s} = \frac{T_{a,s+1}T_{a,s-1}}{T_{a+1,s}T_{a-1,s}} \,. \tag{2.30}$$

 T_{as} should satisfy the Hirota equations

$$T_{a,s}^2 = T_{a+1,s}T_{a-1,s} + T_{a,s+1}T_{a,s-1} , \qquad (2.31)$$

from which all (2.1-2.7) follow. The indices of these T_{as} functions belong to the *T*-shaped lattice (see Fig.2). Another equivalent representation, which follows from (2.31) is

$$1 + Y_{a,s} = \frac{T_{a,s}^2}{T_{a+1,s}T_{a-1,s}} .$$
(2.32)

It is important to notice that the choice of T_{as} is not unique for given Y_{as} , there is a "gauge" freedom

$$T_{a,s}(z) \to g_1(z)[g_2(z)]^a T_{a,s}(z) ,$$
 (2.33)

which leaves Y's unchanged for two arbitrary functions $g_1(z)$ and $g_2(z)$.

Below we will see how the general solution of (2.31) can be constructed for the infinite vertical strip, which is the upper part of the *T*-hook. Then we constrain it by the large *a* asymptotic (2.26) and (2.27), and match with Y_{\odot_s} given by (2.28) (shown as dark gray circles on Fig.2). As a result all *Y*'s are constructed explicitly in Sec.2.3.2 for finite L/g.

2.3.1 Solution of Hirota equation in the vertical strip

We notice that inside the vertical strip of Fig.2 the Hirota equation (2.31) coincides with the recurrent relation for the characters of SU(4) for the representations with rectangular

⁶the -i0 is due to the prescription to go below singularities in the convolutions [19].



Figure 2: T-shaped "fat hook". Small black filled circles correspond to the finite T_{as} function. Big circles show where solution for Y-functions for the strips is applicable.

Young tableaux. It is well known that the characters are given by Schur polynomials for Young tableau $(\lambda_1, \ldots, \lambda_4)$:⁷

$$s_{\lambda}(y) = \frac{\det(y_i^{\lambda_j+4-j})_{1 \le i,j \le 4}}{\det(y_i^{4-j})_{1 \le i,j \le 4}} \,. \tag{2.34}$$

We only need symmetric solutions $T_{a,-s} = T_{a,s}$ which implies that $y_4 = 1/y_1$ and $y_3 = 1/y_2$. For rectangular representation $\lambda_{i \leq s+2} = 0$, $\lambda_{i>s+2} = a$, from (2.34) we get $T_{a,2} = 1$ and

$$T_{a,1} = \frac{y_1 y_2}{(y_1 - y_2)(y_1 y_2 - 1)} \left(\frac{y_1}{y_1^2 - 1} \left(\frac{1}{y_1^{a-2}} - y_1^{a-2} \right) - \frac{y_2}{y_2^2 - 1} \left(\frac{1}{y_2^{a-2}} - y_2^{a-2} \right) \right) (2.35)$$

This solution of (2.31) has only 2 parameters because it is suitable for a finite semi-infinite strip boundary conditions, appropriate for characters. On the other hand, the most general $s \to -s$ symmetric solution should have 4 free parameters. It is clear that we can get one more by shifting *a* in (2.35) by an arbitrary function. What is also true, although not as trivial, is that we can shift independently *a* in y_1^a and in y_2^a so that the most general solution of (2.31) is

$$\begin{split} T_{a,2} &= 1 , \\ T_{a,1} &= \frac{iy_1y_2}{(y_1 - y_2)(y_1y_2 - 1)} \left(\frac{y_1}{y_1^2 - 1} \left(\frac{1}{S_1y_1^a} + S_1y_1^a \right) - \frac{y_2}{y_2^2 - 1} \left(\frac{1}{S_2y_2^a} + S_2y_2^a \right) \right) (2.36) \\ T_{a,0} &= T_{a,1}^2 - T_{a+1,1}T_{a-1,1} , \end{split}$$

all the others T's are given by $T_{a,-s} = T_{a,s}$.

⁷We would like to thank V.Kazakov for discussing this point.

To establish a relation with the previous section we reparameterize our solution in terms of 4 new parameters in the following way

$$y_1 = \sqrt{\epsilon}$$
, $y_2 = \frac{\sqrt{\epsilon}}{T}$, $S_1 = \sqrt{\frac{SU}{T}}$, $S_2 = \sqrt{\frac{TU}{S}}$. (2.37)

Then

$$T_{a,1} = \frac{i\sqrt{T}\epsilon^{-a/2}}{\sqrt{SU}(1-T)(T-\epsilon)} \left(\frac{T - SU\epsilon^{a+2}}{1-\epsilon} - \frac{ST^{a+3} - U\epsilon^{a+2}/T^a}{T^2 - \epsilon}\right)$$
(2.38)

which for small ϵ coincides with (2.23) up to a gauge transformation (2.33) and thus zero ϵ limit is the asymptotic limit we consider in the previous section. We can now easily compute the large a limit of Y_{\bullet_a} and Y_{\triangle_a} :

$$\lim_{a \to \infty} \frac{\log Y_{\Delta_a}(z-i0)}{a} = \log \frac{1}{T} \quad , \quad \lim_{a \to \infty} \frac{\log Y_{\bullet_a}(z-i0)}{a} = \log \epsilon \; . \tag{2.39}$$

From this, using (2.26) and (2.27), we fix two of the four unknown function,

$$T = \frac{f}{\overline{f}} \quad , \quad \epsilon = \Delta f^2 \; . \tag{2.40}$$

So far we have constructed all Y_{\bullet_a} and Y_{\triangle_a} in terms of two yet unknown functions S(z) and U(z) in such way that (2.2) and (2.3) are already satisfied. We have to glue this solution with the solution inside the right and left wings (2.28) parameterized by a third unknown function A(z) and demand the remaining equations (2.4-2.7,2.12) to be satisfied.

2.3.2 Matching wings

In the previous section we construct the solution for the upper wing of Y-system (the light gray dots on Fig.2) and (2.28) gives the general solution for the left and right wings Y_{\bigcirc_s} . We parameterized these functions in terms of three unknown functions S(z), U(z), A(z). We still have to match these two solutions in the different domains and find the two remaining fermionic Y-functions $Y_{\otimes}(z)$ and $Y_{\oplus}(z)$. To fix these five functions we have exactly five remaining equations (2.4-2.7,2.12). Excluding Y_{\otimes} and Y_{\oplus} we get

$$\frac{Y_{\Delta_2}^2(1+1/Y_{\Delta_3})}{1+Y_{\Phi_2}} = \frac{F}{(A-1)^2} , \qquad (2.41)$$

$$\frac{1+1/Y_{\Delta_2}}{1+Y_{\bullet_1}} = \frac{1}{A^2} \left(\frac{(A-1)^2}{F} - 1\right)^2 , \qquad (2.42)$$

$$Y_{\bullet_1}^2(1+1/Y_{\bullet_2}) = \left(\frac{(A-1)^2(F-1)}{(A-1)^2 - F}\right)^2 , \qquad (2.43)$$

where $F \equiv Y_{\otimes}Y_{\oplus}$. F can be expressed in terms of Y_{\bullet_n} from (2.12):

$$F = \frac{1}{f\bar{f}} \prod_{m=1}^{\infty} (1 + Y_{\bullet_m}) .$$
 (2.44)

Notice that the right hand side of (2.41-2.44) depends on A and F only while the left hand side of these equations does not depend on the two unknown functions S, U and the known

functions f, \bar{f} and Δ . These equations are relatively easy to solve perturbatively in Δ . For example for A with $\mathcal{O}(\Delta^4)$ precision we found

$$A \simeq \frac{1}{\bar{f}-1} + \frac{f}{f-1} + \Delta(\bar{f}+f-2) + \Delta^2(\bar{f}^2+f^2-2) + \Delta^3(\bar{f}^3+f^3-2) + \dots$$

We see that the expansion coefficients are very simple. We can easily sum them up to get the exact result

$$A = -\frac{1}{\Delta \bar{f} - 1} + \frac{ff - 1}{(f - 1)(\bar{f} - 1)} - \frac{1}{f\Delta - 1} + \frac{2}{\Delta - 1},$$

$$F = 2 - \frac{2(f - 1)^2(\bar{f} - 1)^2}{f(\Delta - 1)^3 \bar{f}} - \frac{(f - 1)(5f\bar{f} - 3\bar{f} - 3f + 1)(\bar{f} - 1)}{f(\Delta - 1)^2 \bar{f}} - \frac{(\bar{f} + f - 2f\bar{f})^2}{f(\Delta - 1)\bar{f}} - f\bar{f},$$

$$S = \frac{\bar{f}(f - 1)^2(\bar{f}\Delta - 1)^2}{f(\bar{f} - 1)^2(f\Delta - 1)^2},$$

$$U = \frac{(f - 1)^2(\bar{f} - 1)^2}{f^2 \bar{f}^2(f\Delta - 1)^2(\bar{f}\Delta - 1)^2}.$$
(2.45)

The result is quite simple compared to what one may expect to get from the high degree polynomial equations and apparently there should exist some more straightforward way to get this result. It is easy to check that (2.41, 2.42) and (2.43) are indeed satisfied. To check (2.44) one can use (2.32) to get rid of the infinite product

$$F = \frac{1}{f\bar{f}} \prod_{m=1}^{\infty} (1 + Y_{\bullet_m}) = \frac{1}{f\bar{f}} \prod_{m=1}^{\infty} \frac{T_{m,0}^2}{T_{m-1,0}T_{m+1,0}} = \frac{1}{f\bar{f}} \frac{T_{1,0}}{T_{0,0}} \lim_{a \to \infty} \frac{T_{a,0}}{T_{a+1,0}} = \frac{T_{1,0}}{T_{0,0}} \Delta , \quad (2.46)$$

which allows to express the r.h.s. as a rational function of ϵ, U, S and T so that (2.45) can be easily checked. Notice that we found all Y-functions except for the fermionic ones for which we only presented explicitly the form of their product F. However, using e.g. (2.7) we can easily get Y_{\otimes} in terms of the other Y's that we just fixed.

We can now plug the functions we just found into the $Y_{a,s}$ functions to get explicit expressions for all of these functions in terms of f, \bar{f} and Δ alone! We recall that these functions are completely fixed in terms of the Bethe roots (2.11,2.17). Since the results are not particularly simple, we present them in Appendix B in Mathematica form.

2.4 Energy and Momentum

Having all Y's computed we can easily evaluate the energy of the state, corresponding to a given distribution of roots from (1.4). Using it at strong coupling

$$\epsilon_1^{\rm ph}(z_k) = \frac{x_k^2 + 1}{x_k^2 - 1} + \mathcal{O}\left(\frac{1}{g^2}\right) \quad , \quad \epsilon_a^{\rm mir}(z) = -\frac{iaz}{\sqrt{1 - z^2}} \,, \tag{2.47}$$

and applying the trick from the previous section to compute infinite products of $1 + Y_{as}$ used in (2.46) to get

$$e^{\mathcal{M}_0} \equiv \prod_{a=1}^{\infty} (1+Y_{\bullet_a})^a = \lim_{a \to \infty} \frac{(T_{a,0})^{a+1}}{T_{0,0}(T_{a+1,0})^a} = \frac{(f\Delta - 1)^4 (\bar{f}\Delta - 1)^4}{(\Delta - 1)^4 (f\bar{f}\Delta - 1)^2 (f^2\Delta - 1)(\bar{f}^2\Delta - 1)}.$$
(2.48)

From where we get the following stunningly simple expression, accurate to all orders in wrapping

$$E = \sum_{i=1}^{M} \frac{x_i^2 + 1}{x_i^2 - 1} + \int_{-1}^{1} \frac{dz}{2\pi} \frac{z}{\sqrt{1 - z^2}} \partial_z \mathcal{M}_0 .$$
 (2.49)

Using the expressions from Appendix B one can see that each separate term in (2.48) is significantly more complicated then the resulting product! One can easily check (2.48) using explicit expressions for Y_{\bullet_a} from Appendix C and expanding both sides of the equality in powers of Δ .

In (2.49), the integration goes over $z \in (-1, 1)$ because outside these region Y_{\bullet_a} are strongly suppressed (see (2.25)). Notice that the first term is of order $M \sim g \sim \sqrt{\lambda}$ and contains the classical string energy, whereas the second is ~ 1 and should be a part of the one-loop correction. The reader may already suspect that the numerator of (2.48) corresponds to 4 + 4 fermionic fluctuation modes whereas the terms in the denominator correspond to $4 - \text{modes of } S^5$ and $2 + 1 + 1 - \text{modes of } AdS_5$. We make this relation more precise in Sec.3. To separate the classical string energy from the one-loop corrections in the first term of (2.52) we should find the equation determining positions of z_i with 1/gprecision. In the next section we will consider the exact Bethe ansatz equation (1.5) in the strong coupling limit.

The total momentum of the state can be computed similarly to the energy. One should simply replace the expression for the magnon energy ϵ_a in (1.4) by the magnon momentum

$$P = \sum_{i=1}^{M} \frac{x_i}{g(x_i^2 - 1)} + \frac{1}{2g} \int_{-1}^{1} \frac{dz}{2\pi} \frac{1}{\sqrt{1 - z^2}} \partial_z \mathcal{M}_0 .$$
(2.50)

The natural extension of the cyclicity condition (2.13) is

$$P = 2\pi m . (2.51)$$

This is an additional constrain and one should prove its consistency with the other equations. We will assume (2.51) to be satisfied.

2.5 Exact Bethe ansatz equations

In the previous section we saw that the exact energy of a given state in the semi-classical limit is given by

$$E = \sum_{i=1}^{M} \frac{x_i^2 + 1}{x_i^2 - 1} + \int_{-1}^{1} \frac{dz}{2\pi} \frac{z}{\sqrt{1 - z^2}} \partial_z \log \frac{(f\Delta - 1)^4 (\bar{f}\Delta - 1)^4}{(\Delta - 1)^4 (f\bar{f}\Delta - 1)^2 (f^2\Delta - 1)(\bar{f}^2\Delta - 1)} , \quad (2.52)$$

and the functions f, \bar{f}, Δ determined uniquely in terms of the Bethe roots z_i (2.11,2.17). We still need to find the positions of the roots z_i in order to get the exact energy of the state. In this section we shall derive that these Bethe equations, accurate to all orders in wrapping with one-loop precision, read

$$1 = -\left(\frac{x_k^-}{x_k^+}\right)^L \prod_{j=1}^M \frac{x_k^- - x_j^+}{x_k^+ - x_j^-} \frac{1 - 1/(x_k^+ x_j^-)}{1 - 1/(x_k^- x_j^+)} \sigma^2(z_k, z_j) \\ \times \exp\left[-2\int_{-1}^1 \left(r(x_k, z)\mathcal{M}_+ - r(1/x_k, z)\mathcal{M}_- + u(x_k, z)\mathcal{M}_0\right) dz - iP\right], \quad (2.53)$$

where

$$e^{\mathcal{M}_{+}} = \frac{(f\Delta - 1)^{2}}{(f^{2}\Delta - 1)(f\bar{f}\Delta - 1)} , \ e^{\mathcal{M}_{-}} = \frac{(\bar{f}\Delta - 1)^{2}}{(\bar{f}^{2}\Delta - 1)(f\bar{f}\Delta - 1)} ,$$
(2.54)

The factor $\sigma^2(z_k, z_j)$ contains both the leading Arutyunov-Frolov-Staudacher (AFS) [6] and the sub-leading Hernandez-Lopez (HL) phase [39]. Initially the AFS phase was designed to give an agreement with classical theory. Then it was realized that an extra phase is needed in order to get an agreement with the semi-classical one-loop string energies [38]. Basing on the known expressions for the one-loop energies of particular classical solutions [43] this extra phase was found in [39]. However, in [31] it was shown that even with these both dressing factors the Bethe ansatz equations are misses some exponential corrections. In [24, 25] the one-loop compatibility of the asymptotical Bethe ansatz was proven for a generic classical string motion in $AdS_5 \times S^5$. In [24, 25] it was also noticed that in order to get the agreement one should drop some definite exponential in L/g terms. In the next section we will work with finite L/g keeping all the previously dropped terms and show that the above equation, obtained from the Y-system, describes accurately the one-loop string energies for the $\mathfrak{sl}(2)$ sub-sector.

We now derive the above mentioned equations. To find the exact positions of the Bethe roots z_i one has to evaluate $Y_{\bullet_1}(z_i)$ on the physical real axis (see Fig.3). The results we obtain in the previous sections are applicable in the domains where Y's are smooth functions, however close to the Bethe roots one may expect poles and the approximation used so far is no longer valid.

To get round these difficulties we use the representation of the Y-system obtained in [19], based on the TBA approach for the ground state [20, 19]. The equation we need is the integral equation for the middle node

$$\log Y_{\bullet_1} = \mathcal{T}_{1m} * \log(1 + Y_{\bullet_m}) + 2\mathcal{R}^{(10)} \circledast \log(1 + Y_{\otimes}) + 2\mathcal{R}^{(10)} \circledast K_{m-1} * \log(1 + Y_{\Delta_m}) + i\Phi , \quad (2.55)$$

where \mathcal{T}_{1m} is a kernel containing the dressing phase. * stands for convolution with integration over real axis, whereas \circledast is a convolution with integration around the cut (-1, 1)see Appendix A for more details. To use this equation we will need to know Y_{Δ_m} and Y_{\bullet_m} on the whole real axis. In Sec.2.1 we noticed that for |z| > 1 the asymptotic solution from Sec.2.1 can be used since $Y_{\bullet_a}(z)$ are strongly suppressed for these values of z whereas for |z| < 1 the solution of the Y-system was built above.

We denote by Y_{as}^0 the asymptotic solution, constructed for a given set of *exact* Bethe roots z_i in Sec.2.2, satisfying the exact Bethe equation $Y_{\bullet_1}(z_i) = -1$. The asymptotic solution should satisfy the following integral equation

$$\log Y_{\bullet_1}^0 = 2\mathcal{R}^{(10)} \otimes \log(1 + Y_{\otimes}^0) + 2\mathcal{R}^{(10)} \otimes K_{m-1} * \log(1 + Y_{\Delta_m}^0) + i\Phi , \qquad (2.56)$$



Figure 3: Branch-cut structure of Y_{\bullet_1} . Dashed line is a "mirror" real axis. Between the branch cuts all Y_{\bullet_a} are exponentially suppressed at strong coupling and one can use the asymptotic solution here. To find the exact positions of the Bethe roots one should analytically continue Y_{\bullet_1} under the upper cut to reach a "physical" real axis.

we can subtract it from (2.55) to get

$$\log \frac{Y_{\bullet_1}}{Y_{\bullet_1}^0} = \mathcal{T}_{1m} * \log(1+Y_{\bullet_m}) + 2\mathcal{R}^{(10)} \circledast \log\left(\frac{1+Y_{\otimes}}{1+Y_{\otimes}^0}\right) + 2\mathcal{R}^{(10)} \circledast K_{m-1} * \log\left(\frac{1+Y_{\Delta_m}}{1+Y_{\Delta_m}^0}\right) = \mathcal{T}_{1m} * \log(1+Y_{\bullet_m}) + 2\mathcal{R}^{(10)} \circledast \log\left(\frac{1+Y_{\otimes}}{1+Y_{\otimes}^0}\right) + 2\mathcal{R}^{(10)} \circledast K_{m-1} * \log\left(\frac{1+Y_{\Delta_m}}{1+Y_{\Delta_m}^0}\right) = \mathcal{T}_{1m} * \log(1+Y_{\bullet_m}) + 2\mathcal{R}^{(10)} \circledast \log\left(\frac{1+Y_{\otimes}}{1+Y_{\otimes}^0}\right) + 2\mathcal{R}^{(10)} \circledast K_{m-1} * \log\left(\frac{1+Y_{\Delta_m}}{1+Y_{\Delta_m}^0}\right) = \mathcal{T}_{1m} * \log(1+Y_{\bullet_m}) + 2\mathcal{R}^{(10)} \circledast \log\left(\frac{1+Y_{\otimes}}{1+Y_{\otimes}^0}\right) + 2\mathcal{R}^{(10)} \circledast K_{m-1} * \log\left(\frac{1+Y_{\Delta_m}}{1+Y_{\Delta_m}^0}\right) = \mathcal{T}_{1m} * \log(1+Y_{\bullet_m}) + 2\mathcal{R}^{(10)} \circledast \log\left(\frac{1+Y_{\otimes}}{1+Y_{\otimes}^0}\right) + 2\mathcal{R}^{(10)} \circledast K_{m-1} * \log\left(\frac{1+Y_{\Delta_m}}{1+Y_{\Delta_m}^0}\right) = \mathcal{T}_{1m} * \log\left(\frac{1+Y_{\Delta$$

As a result of these trick we do not need any more to go outside -1 < z < 1 region in the convolutions, since the integrands vanish there. We can now analytically continue $Y_{\bullet_1}/Y_{\bullet_1}^0$ to the physical real axis where the Bethe roots are situated (see Fig.3). A similar analytical continuation was already performed in [8]

$$\log \frac{Y_{\bullet_{1}}^{\text{ph}}}{Y_{\bullet_{1}}^{\text{ph0}}} = \mathcal{T}_{1m}^{\text{ph,mir}} * \log(1+Y_{\bullet_{m}}) + 2\mathcal{R}^{(10)\text{ph,mir}} \otimes \log\left(\frac{1+Y_{\otimes}}{1+Y_{\otimes}^{0}}\right)$$
(2.57)
+ $2\mathcal{R}^{(10)\text{ph,mir}} \otimes K_{m-1} * \log\left(\frac{1+Y_{\Delta_{m}}}{1+Y_{\Delta_{m}}^{0}}\right) + 2K_{m-1}(z_{k}-\frac{i}{4g}) * \log\left(\frac{1+Y_{\Delta_{m}}}{1+Y_{\Delta_{m}}^{0}}\right) .$

Now we simply have to expand the kernels at large g and substitute Y's. To expand $\mathcal{T}_{1m}^{\text{ph,mir}}$ again can use the AFS dressing phase (2.24)

$$\mathcal{R}^{(10)\text{ph,mir}}(z_k, w) \simeq r(x_k, w) , \quad \mathcal{B}^{(10)\text{ph,mir}}(z_k, w) \simeq r(1/x_k, w) , \mathcal{T}_{1m}^{\text{ph,mir}}(z_k, w) \simeq -m \left[2r(x_k, w) + 2u(x_k, w) + p(w)\right] , \qquad (2.58) K_m(z_k - w) \simeq \delta(w - z_k) + m \left[r(x_k, w) + r(1/x_k, w)\right] ,$$

where we use the following notations

$$r(x,z) = \frac{x^2}{x^2 - 1} \frac{\partial_z}{2\pi g} \frac{1}{x - x(z)} \quad , \quad u(x,z) = \frac{x}{x^2 - 1} \frac{\partial_z}{2\pi g} \frac{1}{x^2(z) - 1} \quad , \quad p(z) = \frac{\partial_z}{4\pi i g} \frac{1}{\sqrt{1 - z^2}} \quad .$$

After that we can rearrange the terms in (2.57) according to kernels r. Using the following



Figure 4: Density as a function of z for a typical distribution of bethe roots z_i on the real axis.

"magic" products⁸

$$e^{-\mathcal{M}_{+}} \equiv \frac{1+Y_{\otimes}}{1+Y_{\otimes}^{0}} \prod_{m=2}^{\infty} \left(\frac{1+Y_{\Delta_{m}}}{1+Y_{\Delta_{m}}^{0}}\right)^{m} \prod_{m=1}^{\infty} \frac{1}{(1+Y_{\bullet_{m}})^{m}} = \frac{(f^{2}\Delta-1)(f\bar{f}\Delta-1)}{(f\Delta-1)^{2}} \quad (2.59)$$
$$e^{+\mathcal{M}_{-}} \equiv \frac{1+1/Y_{\oplus}^{0}}{1+1/Y_{\oplus}} \prod_{m=2}^{\infty} \left(\frac{1+Y_{\Delta_{m}}}{1+Y_{\Delta_{m}}^{0}}\right)^{m-2} = \frac{(\bar{f}\Delta-1)^{2}}{(\bar{f}^{2}\Delta-1)(f\bar{f}\Delta-1)} , \quad (2.60)$$

we get the corrected ABA equations (2.53) accurate to all orders in wrapping with one-loop precision. Here we use the expression for the total momentum (2.50). Notice that the last term in the exponent is irrelevant due to (2.51).

2.6 Finite gap solutions

In this section we expand the corrected Bethe ansatz equation (2.53), obtained in the previous section for a particular type of configurations of roots. Before expanding (2.53) one should take log of both sides. Due to the $2\pi i$ ambiguity of the log function one should assign an integer mode number n_k for each root x_k , i.e. for each of the M equations we can assign $\log(1) = 2\pi i n_k$ in the left hand side of (2.53). For the finite gap solutions we assume that the set of mod numbers n_k contains only a finite number of different integers. In this limit the Bethe roots x_k are forming dense distributions along some cuts C_n on the real axis with some density $\rho(x)$ so that the number of cuts is equal to the number of distinct mod numbers n_k (see Fig.4). One can establish a one-to-one correspondence between such configurations and finite gap classical strings motions considered in the next section.

Strong coupling expansion of ABA for these configuration was studied intensively [37, 38, 23, 39]. Here we can use the existing results to expand our corrected by wrappings Bethe equation (2.53). Usually one defines the quasi-momenta through resolvent (2.10)

$$p_{\hat{2}}(x) = \frac{\frac{L}{2g}x + G(0)}{x^2 - 1} + G(x) .$$
(2.61)

⁸To compute these products we again use z-i0 prescription to ensure their convergence. This prescription is inherited from the TBA equation where the integration should go slightly below the real axis.

From the definition of the resolvent (2.10) we see that p(x) has poles at $x = x_k$. When the number of roots goes to infinity these poles condense into branch cuts and we can rewrite the quasi-momenta in terms of the density ρ of the roots

$$p_{\hat{2}}(x) = \frac{\frac{L}{2g}x + G(0)}{x^2 - 1} + \sum_{n} \int_{\mathcal{C}_n} \frac{\rho(y)}{x - y} \frac{y^2}{y^2 - 1} dy .$$
 (2.62)

Equation (2.53) gives an integral equation on the density of the roots

$$2\pi n = \left[p_{\hat{2}}(x+i0) + p_{\hat{2}}(x-i0) \right] + \alpha(x) p_{\hat{2}}'(x) \cot p_{\hat{2}}(x) + \mathcal{V}(x) - 2i \int_{-1}^{1} \left(r(x,z)\mathcal{M}_{+} - r(1/x,z)\mathcal{M}_{-} + u(x,z)\mathcal{M}_{0} \right) dz , \quad x \in \mathcal{C}_{n} , \qquad (2.63)$$

where $\alpha(x) = \frac{x^2}{g(x^2-1)}$. The second term in (2.63) is so-called "anomaly" term [37, 38, 41, 23] and the third therm is the contribution of subleading term in the dressing phase – the Hernandez-Lopez phase [39]

$$\mathcal{V}(x) = \alpha(x) \sum_{\substack{r,s=2\\r+s \in Odd}}^{\infty} \frac{1}{\pi} \frac{(r-1)(s-1)}{(s-r)(r+s-2)} \left(\frac{Q_r}{x^s} - \frac{Q_s}{x^r}\right) , \quad Q_s = \oint \frac{dx}{2\pi i} \frac{G(x)}{x^s} . \tag{2.64}$$

The last line in (2.63) incorporates the finite size effects. Finally using the standard notations

$$p_{\hat{1}}(x) = -p_{\hat{2}}(1/x) , \quad p_{\hat{2}}(x) = \frac{\frac{L}{2g}x + G(0)}{x^2 - 1}$$
 (2.65)

we can rewrite \mathcal{M} 's in terms of the quasi-momenta

$$e^{\mathcal{M}_{0}} = \frac{(1 - e^{-ip_{\hat{2}} - ip_{\hat{2}}})^{4}(1 - e^{-ip_{\hat{1}} - ip_{\hat{2}}})^{4}}{(1 - e^{-2ip_{\hat{2}}})^{4}(1 - e^{-ip_{\hat{1}} - ip_{\hat{2}}})^{2}(1 - e^{-2ip_{\hat{2}}})(1 - e^{-2ip_{\hat{1}}})},$$

$$e^{\mathcal{M}_{+}} = \frac{(1 - e^{-ip_{\hat{2}} - ip_{\hat{2}}})^{2}}{(1 - e^{-2ip_{\hat{2}}})(1 - e^{-ip_{\hat{1}} - ip_{\hat{2}}})},$$

$$e^{\mathcal{M}_{-}} = \frac{(1 - e^{-ip_{\hat{1}} - ip_{\hat{2}}})^{2}}{(1 - e^{-2ip_{\hat{1}}})(1 - e^{-ip_{\hat{1}} - ip_{\hat{2}}})}.$$
(2.66)

In the next section we will see how these structures appear in the quasi-classical string quantization.

We have found a set of equations which are supposed to correct the Beisert-Staudacher asymptotic equations with the Beisert-Eden-Staudacher dressing phase in the strong coupling scaling limit. The latter are known to describe the semi-classical string spectrum up to exponentially suppressed finite size corrections as described in the previous section. This extra corrections which we just derived ought to cure the known mismatch and correctly incorporate *all* wrapping corrections to 1-loop precision. In the next section we show that this turns out to be precisely the case!

3. Quasi-classical string quantization

In this section we review the quasi-classical quantization method. Then we consider a generic solution of string equations of motion inside $AdS_3 \times S^1$ and compute its one-loop energy.

The one-loop correction to a classical string energy could be understood as zero point oscillations of fluctuations around the classical solution. To compute it one can expand the classical action up to the quadratic order around the classical solution and then find the spectrum of oscillation modes. These modes could be labeled by the mode number n, which tells us how many wavelengthes fit into the string, and polarization. There are 8+8 bosonic and fermionic polarizations which we label by double indices (ij):

Bosonic :
$$(\hat{1},\hat{3}), (\hat{1},\hat{4}), (\hat{2},\hat{3}), (\hat{2},\hat{4}), (\tilde{1},\tilde{3}), (\tilde{1},\tilde{4}), (\tilde{2},\tilde{3}), (\tilde{2},\tilde{4}),$$

Fermionic : $(\hat{1},\tilde{3}), (\hat{1},\tilde{4}), (\hat{2},\tilde{3}), (\hat{2},\tilde{4}), (\tilde{1},\hat{3}), (\tilde{1},\hat{4}), (\tilde{2},\hat{3}), (\tilde{2},\hat{4}).$ (3.1)

We denote the energies of the vibrations Ω_n^{ij} . Then the one-loop correction is simply a sum of halves of these fluctuation energies [40, 46]

$$\delta E_{1-loop} = \frac{1}{2} \sum_{n,(ij)} (-1)^{F_{ij}} \Omega_n^{ij} , \qquad (3.2)$$

where F_{ij} is +1 for bosonic polarizations and -1 for fermionic.

The direct computation of these $\Omega_n^{(ij)}$ is only possible in the simplest cases [42, 43]. For a generic solution it is enormously hard to perform this calculation starting from the classical action. The tool which allows to handle the quasi-classical string quantization efficiently is the algebraic curve technique developed in [23, 26, 27, 29]. Below we describe the construction of the algebraic curve and the method of the quasi-classical calculations.

3.1 Classical algebraic curve

The classical equations of motion of the Metsaev-Tseytlin superstring action [44] can be summarized in a compact form as the flatness condition [4]

$$dA - A \wedge A = 0. \tag{3.3}$$

for a connection $A(\sigma, \tau; x)$ which is a local functional of fields depending on an arbitrary complex number called the spectral parameter x and taking its values in $\mathfrak{psu}(2,2|4)$. The fact that the classical equations of motion could be packed into the flatness condition is an indication that the model is classically integrable. Indeed, we can define the monodromy matrix

$$M(x) = \operatorname{Pexp} \oint_{\gamma} A(x) \tag{3.4}$$

where γ is a loop wrapping the worldsheet cylinder once. The flatness of the connection ensures path independence of the spectral data of the super $(4 + 4) \times (4 + 4)$ matrix M(x). In particular, the displacement of the whole loop in time direction amounts to a similarity transformation and we conclude that the eigenvalues of the monodromy matrix are conserved with time quantities depending on the spectral parameter x. We denote the eigenvalues of M(x) as

$$\{e^{ip_{\hat{1}}}, e^{ip_{\hat{2}}}, e^{ip_{\hat{3}}}, e^{ip_{\hat{4}}} | e^{ip_{\tilde{1}}}, e^{ip_{\tilde{2}}}, e^{ip_{\tilde{3}}}, e^{ip_{\tilde{4}}}\},$$
(3.5)

where $p_i(x)$ and $p_i(x)$ are so-called quasi-momenta [5, 7]. The quasi-momenta contain information about all conserved charges of the theory, in particular the global symmetry charges, including the energy E. The eigenvalues are the roots of the characteristic polynomial and thus they define an 8-sheet Riemann surface. In general these sheets are connected by several branch-cuts. The branch points on this surface are the values of the spectral parameter x where two eigenvalues coincide and M(x) cannot be diagonalized completely. Different classical solutions correspond to different algebraic curves. For many calculation the explicit construction of the classical solution in terms of the initial fields entering into the Lagrangian is not needed and can be replaced by the corresponding algebraic curve. For example the energy can be computed as a simple contour integral

$$E = \sum_{\mathcal{C}_{ij}} 2g \oint_{\mathcal{C}_{ij}} \frac{dx}{2\pi i} \frac{p_i}{x^2} \,. \tag{3.6}$$

It is always possible to define the quasi-momenta so that they vanish at $x \to \infty$. Then, however, the quasi-momenta should jump by a multiple of 2π when passing through a cut

$$\frac{p_i(x+i0) + p_i(x-i0)}{2} - \frac{p_j(x+i0) + p_j(x-i0)}{2} = 2\pi n \quad , \quad x \in \mathcal{C} \; . \tag{3.7}$$

The quasi-momenta are restricted by the properties of the monodromy matrix M(x). Due to super-tracelessness

$$p_{\hat{1}} + p_{\hat{2}} + p_{\hat{3}} + p_{\hat{4}} = p_{\tilde{1}} + p_{\tilde{2}} + p_{\tilde{3}} + p_{\tilde{4}} , \qquad (3.8)$$

and as a consequence of the special properties of M(x) under $x \to 1/x$ transformation one has

$$p_{\hat{1}}(x) = -p_{\hat{2}}(1/x) , \quad p_{\tilde{1}}(x) = -2\pi m - p_{\tilde{2}}(1/x)$$

$$p_{\hat{4}}(x) = -p_{\hat{3}}(1/x) , \quad p_{\tilde{4}}(x) = +2\pi m - p_{\tilde{3}}(1/x) ,$$
(3.9)

where m is an integer winding number.

There are also infinitely many points where two eigenvalues coincide but, nevertheless, the matrix M(x) can be diagonalized. The two quasi-momenta p_i and p_j corresponding to the coincident eigenvalues have no singularity and differs by $2\pi n$. One can perturb the curve by opening a small cut connecting the intersecting sheets of the surface at these points. We label these points by an integer n and a couple of indices (ij)

$$p_i(x_n^{ij}) - p_j(x_n^{ij}) = 2\pi n . (3.10)$$

One of the nice features of the algebraic curve is the simplicity of visualization of the action variables of this classical integrable theory. They are the contour integrals around the branch cuts

$$\frac{g}{2\pi i} \oint_{\mathcal{C}} \left(1 - \frac{1}{x^2} \right) p_i(x) dx .$$
(3.11)

In the standard quasi-classical quantization procedure one should assume them to be integers.

3.1.1 Algebraic curve for $\mathfrak{sl}(2)$ subsector

The algebraic curve for the string in $AdS_3 \times S^1$ was constructed in [45]. In the general framework reviewed in the previous section this sector corresponds to the cuts connecting $p_{\hat{2}}$ with $p_{\hat{3}}$ outside the unit circle centered at the origin. Automatically, due to the $x \to 1/x$ symmetry (3.9), we will have reflected cuts connecting $p_{\hat{1}}$ with $p_{\hat{4}}$ inside the unit circle. One can easily build the spectral representation for the quasi-momenta [45]

$$p_{\hat{2}}(x) = -p_{\hat{3}}(x) = -p_{\hat{1}}(1/x) = p_{\hat{4}}(1/x) = \frac{\frac{L}{2g}x + G(0)}{x^2 - 1} + G(x) , \qquad (3.12)$$

$$p_{\tilde{2}}(x) = -p_{\tilde{3}}(x) = p_{\tilde{1}}(x) = -p_{\tilde{4}}(x) = \frac{\frac{1}{2g}x + G(0)}{x^2 - 1}, \qquad (3.13)$$

where $G(x) = \int_{\mathcal{C}} \frac{\rho(y)}{x-y} \frac{y^2}{y^2-1} dy$. We see that these quasi-momenta are exactly those of Sec.2.6. The action variables (3.11) count the number of Bethe roots constituting the cut. In this way one established the map between classical solutions and the finite gap configurations of Bethe roots [45].

The equation (3.7) for the $\mathfrak{sl}(2)$ subsector becomes

$$p_{\hat{2}}(x+i0) + p_{\hat{2}}(x-i0) = 2\pi n , \ x \in \mathcal{C} ,$$
 (3.14)

which is now an integral equation for the discontinuity $\rho(x)$.

3.2 Quasi-classical corrections from algebraic curve

Using the algebraic curve it is also possible to find the spectrum of the fluctuations $\Omega_n^{(ij)}$ around an arbitrary classical solution using the techniques developed in [23, 24, 25, 26, 27, 29]. The perturbations of the given classical solution are reflected in the algebraic curve by extra cuts. The small cuts could only appear in the special points of the curve given by (3.10). The perturbed quasi-momenta differ from the non-perturbed ones by a small amount $\delta p_i(x)$. The minimal size of the cut is restricted in the quasi-classically quantized theory by the condition that the contour integral around this new cut (3.11) is integer.

From far away the branch points almost merge and the cut looks like a pole with a tiny residue

$$\delta p_i(x) \sim \frac{\alpha(x)}{x - x_n^{ij}} , \ \alpha(x) = \frac{x^2}{g(x^2 - 1)} ,$$
 (3.15)

such that (3.11) counts a single quantum. We see that for given n and (ij) the perturbation of the quasi-momenta is pretty much restricted and one can compute the energy shift due to this fluctuation. This gives precisely Ω_n^{ij} . In all details this technique is described in [26, 27, 29, 30] (see also [46]).

 Ω_n^{ij} has two contributions different by their nature. Firstly, the extra small cut carries its own energy as we can see from (3.6)

$$\Omega_n^{0\,ij} = \omega(x_n^{ij}) \ , \ \omega(x) \equiv \frac{2}{x^2 - 1} \ ,$$
(3.16)



Figure 5: Complex plane of the spectral parameter x. Positions of the fluctuations x_n^{ij} are shown by the gray dots.

secondly, it deforms others cuts, changing thus their contribution. This second contribution we study in the next section. Let us now see the effect of (3.16) on the one-loop shift (3.2). We have to compute the following sum

$$\delta E_{1-loop}^{0} = \sum_{n} \sum_{(ij)} (-1)^{F_{ij}} \omega \left(x_{n}^{ij} \right) , \qquad (3.17)$$

where x_n^{ij} should be found from (3.10). We rewrite this sum over n as an integral

$$\delta E_{1-loop}^{0} = \sum_{(ij)} (-1)^{F_{ij}} \oint \frac{dn}{4i} \cot(\pi n) \omega\left(x_n^{ij}\right) , \qquad (3.18)$$

where the contour encircles the real axis. Next, for each polarization (ij) we change the integration variable from n to x via (3.10). The integration over n maps to contours which encircle the fluctuation positions x_n^{ij} located outside the unit circle U. Then we can deform this contour in the x plane to get an integral over the unit circle, centered at the origin (see Fig.5)⁹

$$\delta E_{1-loop}^{0} = \sum_{(ij)} (-1)^{F_{ij}} \oint_{\mathbb{U}} \frac{dx}{4i} \frac{p'_i - p'_j}{2\pi} \cot\left(\frac{p_i - p_j}{2}\right) \omega\left(x\right) = \oint_{\mathbb{U}^+} \frac{dx}{2\pi i} \omega\left(x\right) \partial_x \mathcal{N}_0 , \quad (3.19)$$

where \mathbb{U}^+ is the upper half of the unit circle and

$$e^{\mathcal{N}_0} = \prod_{(ij)} \left(1 - e^{-ip_i + ip_j} \right)^{F_{ij}} .$$
(3.20)

We use that

$$\sum_{(ij)} (-1)^{F_{ij}} (p_i - p_j) = 0.$$
(3.21)

⁹For each (ij) the sum in n in (3.17) is divergent and one should carefully treat the large n's. The oneloop shift written as an integral (3.19) corresponds to a particular prescription of the large n regularization which is analyzed in details in [25]. Whereas from the algebraic curve point of view this prescription is absolutely reasonable from the worldsheet action point of view the complete prove is still missing. We should notice however that this question is not a particularity of the finite size effects, it is rather related to the question of the validity of the ABA at the one-loop level.

The product in (3.20) goes over all 8 + 8 polarizations listed in (3.1). Notice that from (3.12) $p_i - p_j \sim L/g$ and the integral (3.20) is exponentially suppressed for L/g large. This kind of terms are not captured by the ABA and as a result the ABA can be only used when L/g is sufficiently large.

To our deep satisfaction we notice that for the $\mathfrak{sl}(2)$ subsector $\mathcal{N}_0 = -\mathcal{M}_0$ from (2.66)! Moreover by changing the integration variable to $z = \frac{1}{2}(x + 1/x)$ we map the integration contour to [-1, 1] segment of the real axis and (3.19) matches precisely with the second term of the expression for the energy obtained from Y-system (2.52)! In the next section we show how the corrected Bethe equation (2.63) arises from the quasi-classical quantization.

3.3 Back-reaction

So far only the direct contribution of the virtual sea of the fluctuations was computed. We have to take into account the back-reaction – the deformation of the quasi-momenta close to the cuts of the initial non-perturbed classical curve. In [25] such deformations were considered dropping exponentially suppressed finite size corrections. This allowed for a precise derivation of the HL correction to the AFS asymptotic Bethe equations. Here we will keep all exponentially suppressed terms since we want to derive a set of exact integral equations. We split $p_{\hat{2}}$ into the part containing all the small virtual cuts $V_{\hat{2}}$ and the smooth part $p_{\hat{2}}^{\text{br}}$. To write down $V_{\hat{2}}$ one should take into account $x \to 1/x$ symmetry and some further analyticity constraints such as poles at $x = \pm 1$. The basic rule is that each fluctuation Ω_n^{ij} contributes as a pole at $1/x_n^{ij}$ due to the constraint (3.9) (see [25] for more details)¹⁰

$$V_{\hat{2}} = \sum_{n} \left(\frac{\alpha(x)}{x} \sum_{(ij)} \frac{(-1)^{F_{ij}}}{(x_n^{ij})^2 - 1} + \alpha(x) \sum_{(\hat{2}j)} \frac{(-1)^{F_{ij}}}{x - x_n^{ij}} - \alpha(1/x) \sum_{(\hat{1}j)} \frac{(-1)^{F_{ij}}}{1/x - x_n^{ij}} \right) .$$
(3.22)

The second sum goes over all fluctuations starting at p_2

Bosonic : $(\hat{2}, \hat{3}), (\hat{2}, \hat{4}),$ Fermionic : $(\hat{2}, \tilde{3}), (\hat{2}, \tilde{4}),$

and in the last term, corresponding to the reflected poles, the sum goes over all fluctuations starting at $p_{\hat{1}}$.

Now we should use (3.14) to find the discontinuity of $p_{\hat{2}}^{\text{br}}$

$$p_{\hat{2}}^{\text{br}}(x+i0) + p_{\hat{2}}^{\text{br}}(x-i0) + 2V_{\hat{2}}(x) = 2\pi n , \ y \in \mathcal{C} .$$
 (3.23)

We can convert the sum over n in (3.22) into the integral over x – precisely like we did with the energy and then deform the contour to the unit circle

$$V_{\hat{2}} = \oint_{\mathbb{U}^+} \frac{dy}{2\pi i} \left(\frac{\alpha(x)}{x} \frac{\partial_y \mathcal{N}_0}{y^2 - 1} + \alpha(x) \frac{\partial_y \mathcal{N}_+}{x - y} - \alpha(1/x) \frac{\partial_y \mathcal{N}_-}{1/x - y} \right) + \frac{\alpha(x)}{2} p_{\hat{2}}' \cot p_{\hat{2}} + \frac{1}{2} \mathcal{V}(x).$$
(3.24)

 $^{^{10}}$ compare with equation 20 in [25].

There is one important difference that now there is an extra pole at x = y caught when deforming the contour to the unit circle given rise to the second term (see [24]). We denote

$$e^{\mathcal{N}_{+}} = \prod_{(\hat{2}j)} \left(1 - e^{-ip_{\hat{2}} + ip_{j}} \right)^{F_{\hat{2}j}} , \quad e^{\mathcal{N}_{-}} = \prod_{(\hat{1}j)} \left(1 - e^{-ip_{\hat{1}} + ip_{j}} \right)^{F_{\hat{1}j}}$$
(3.25)

and

. .

$$\mathcal{V}(x) = \oint_{\mathbb{U}^+} \frac{dy}{2\pi} \left(\frac{\alpha(x)}{x-y} - \frac{\alpha(1/x)}{1/x-y} \right) \partial_y(p_{\hat{1}} + p_{\hat{2}} - p_{\tilde{1}} - p_{\tilde{2}}) .$$
(3.26)

taking into account that $\partial_y(p_{\hat{1}} + p_{\hat{2}} - p_{\tilde{1}} - p_{\hat{2}}) = \partial_y(G(y) - G(1/y))$ one can see that this coincides precisely with the contribution of the Hernandez-Lopez phase [25] in the Y-system analysis (2.64). For $\mathfrak{sl}(2)$ we again have $\mathcal{N}_{\pm} = -\mathcal{M}_{\pm}$ and after change of the integration variable to $z = \frac{1}{2}(x + 1/x)$ we get precisely the equation obtained in the Ysystem framework (2.63). Thus we established the match of these two completely different approaches at the level of equations.

4. Summary and future directions

In this paper we studied the finite size effects at strong coupling for strings in $AdS^3 \times S_1$. We attacked the problem from two directions – from the quasi-classical string quantization using the algebraic curve techniques [26, 27, 29, 30] and from the recently conjectured Y-system [9]. We found the same result in both cases thus providing a very nontrivial test of the latter. We also derived the corrected expression for the energy of (2.52)

$$E = \sum_{i=1}^{M} \frac{x_i^2 + 1}{x_i^2 - 1} + \int_{-1}^{1} \frac{dz}{2\pi} \frac{z}{\sqrt{1 - z^2}} \partial_z \mathcal{M}_0 \quad , \quad e^{\mathcal{M}_0} = \frac{(f\Delta - 1)^4 (\bar{f}\Delta - 1)^4}{(\Delta - 1)^4 (\bar{f}\bar{f}\Delta - 1)^2 (f^2\Delta - 1)(\bar{f}^2\Delta - 1)} \quad .$$

$$\tag{4.1}$$

where f and \bar{f} (2.11) are some simple functions of the Bethe roots x_i and Δ (2.17) is the exponential wrapping parameter. The last integral term is responsible for the finite size effects and vanishes in the large volume limit. We also found that the Bethe roots should satisfy the corrected Bethe equation

$$-1 = \left(\frac{x_k^-}{x_k^+}\right)^L \prod_{j=1}^M \frac{x_k^- - x_j^+}{x_k^+ - x_j^-} \frac{1 - 1/(x_k^+ x_j^-)}{1 - 1/(x_k^- x_j^+)} \sigma^2(z_k, z_j) \\ \times \exp\left[-2 \int_{-1}^1 \left(r(x_k, z)\mathcal{M}_+ - r(1/x_k, z)\mathcal{M}_- + u(x_k, z)\mathcal{M}_0\right) dz\right], \quad (4.2)$$

where \mathcal{M}_+ and \mathcal{M}_- contain all exponential wrapping corrections

$$e^{\mathcal{M}_{+}} = \frac{(f\Delta - 1)^{2}}{(f^{2}\Delta - 1)(f\bar{f}\Delta - 1)} , \quad e^{\mathcal{M}_{-}} = \frac{(\bar{f}\Delta - 1)^{2}}{(\bar{f}^{2}\Delta - 1)(f\bar{f}\Delta - 1)} .$$
(4.3)

There are many interesting directions which would be worth exploring:

• It would be interesting to make a more direct analysis by solving our corrected equations (4.2) for some simple configuration of roots and comparing the solution with the sum of fluctuation energies for the corresponding classical solution. For example, it would be very nice to repeat the analysis of [31] for circular strings in AdS_3 using the corrected Bethe equations.

- It would be also interesting to compare the corrected Bethe equations (4.2) with the conjectured generalized Lüscher formula [47] in the strong coupling scaling limit.
- In this paper we focused on strings moving in $AdS^3 \times S_1$. From the Y-system point of view this is an important simplification because the excited states integral equations are only available for this sector[19]. On the other hand, from the string semi-classics point of view, following [24, 25], the derivation of the corrected Bethe equations would be a straightforward task. It would be very interesting to perform this generalization and to use it as a guiding principle to construct the Y-system integral equations for any excited state.
- It would also be very important to consider analytically some states which cannot be treated by the scaling limit like Konishi state.
- The Y-system conjectured in [9] for Aharony-Bergman-Jafferis-Maldacena [50] theory recently was supported from the 4-loop perturbation theory [51] at weak coupling. It would be also interesting to make some strong coupling test of this conjecture.
- Related to finite size corrections but at weak coupling one should reproduce the results of [48] from the Y-system set of equations.
- Finally one can try to generalize the approach used here to solve the Y-system at finite coupling by bringing it to a couple of integral equations like in [16] (see [49] for some first steps).

In short, there are many interesting open problems to address related to the exact computation of the AdS/CFT planar spectrum and many simplifications are to be expected. We are getting closer and closer to finding the exact solution to a four dimensional superconformal gauge theory for the very first time. The methods developed here could be also useful for a wide range of integrable theories. The quasi-classical quantization probes the theories at finite volume and provides an important information about hidden structures, such as Y-systems.

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A. Notations

There are two distinct possibilities to define x(z) which is solution to x + 1/x = 2z

$$x^{\text{ph}}(z) \equiv z + \sqrt{z - 1}\sqrt{z + 1} , \ x^{\text{mir}}(z) \equiv z + i\sqrt{1 - z^2} .$$
 (A.1)

By default we always choose $x = x^{\text{mir}}$. These two functions coincide above the real axis and have the following properties under complex conjugation

$$\overline{x^{\mathrm{ph}}} = x^{\mathrm{ph}}$$
, $\overline{x} = 1/x$. (A.2)

We also use the notation for the Bethe roots

$$x_j = x^{\text{ph}}(z_j) , \ x_j^{\pm} = x^{\text{ph}}(z_j \pm \frac{i}{4g}) .$$
 (A.3)

The single magnon energy and momentum are

$$\epsilon_a(z) = a + \frac{2ig}{x(z + \frac{ia}{4g})} - \frac{2ig}{x(z - \frac{ia}{4g})} , \quad \pi_a(z) = \frac{1}{i} \log \frac{x(z + \frac{ia}{4g})}{x(z - \frac{ia}{4g})} , \quad (A.4)$$

depending on which x(z) we are using it could be denoted $\epsilon_a^{\rm ph}(z)$ or $\epsilon_a^{\rm mir}(z)$.

The kernels we are using in the integral equations are defined as follows¹¹

$$\begin{split} K_{n}(z) &\equiv \frac{4gn}{\pi(n^{2} + 16g^{2}z^{2})} ,\\ K_{nm}(z) &\equiv \sum_{k_{1}=\frac{1-n}{2}}^{\frac{n-1}{2}} \sum_{k_{2}=\frac{1-m}{2}}^{\frac{m-1}{2}} K_{2}\left(z + i\frac{k_{1}+k_{2}}{2g}\right) ,\\ K_{nm}^{\neq}(z) &\equiv \sum_{k_{1}=\frac{1-n}{2}}^{\frac{n-1}{2}} \sum_{k_{2}=\frac{1-m}{2}}^{\frac{m-1}{2}} K_{1}\left(z + i\frac{k_{1}+k_{2}}{2g}\right) ,\\ \mathcal{R}^{(nm)}(z,w) &\equiv \frac{\partial_{w}}{2\pi i} \log \frac{x(z + \frac{in}{4g}) - x(w - \frac{im}{4g})}{x(z - \frac{in}{4g}) - x(w + \frac{im}{4g})} + \frac{\partial_{w}}{4\pi i} \log \frac{x(w + \frac{im}{4g})}{x(w - \frac{im}{4g})} ,\\ \mathcal{B}^{(nm)}(z,w) &\equiv \frac{\partial_{w}}{2\pi i} \log \frac{1/x(z + \frac{in}{4g}) - x(w - \frac{im}{4g})}{1/x(z - \frac{in}{4g}) - x(w + \frac{im}{4g})} + \frac{\partial_{w}}{4\pi i} \log \frac{x(w + \frac{im}{4g})}{x(w - \frac{im}{4g})} ,\\ \mathcal{M}_{nm} &\equiv K_{n-1} \circledast \mathcal{R}^{(0m)} + K_{n-1,m-1}^{\neq} ,\\ \mathcal{N}_{nm} &\equiv \mathcal{R}^{(n0)} \circledast K_{m-1} + K_{n-1,m-1}^{\neq} , \end{split}$$
(A.6)

There are two types of convolutions * and \circledast . The first corresponds to the usual integration along whole real axis whereas the second one is a convolution a long a path going from -1 to 1 and then back on another sheet e.g.

$$\mathcal{R}^{(n0)} \circledast \log(1+Y_{\otimes}) = \int_{-1}^{1} dz \left[\mathcal{R}^{(n0)} \log(1+Y_{\otimes}) - \mathcal{B}^{(n0)} \log(1+1/Y_{\oplus}) \right]$$

¹¹We use rescaled kernels compare to [19]: $\mathcal{K}^{\text{new}}(z, w) = 2g\mathcal{K}^{\text{old}}(2gz, 2gw).$

where $1/Y_{\oplus}$ is the analytical continuation of Y_{\otimes} across the cut $u \in (-\infty, -1) \cup (1, +\infty)$. The kernel $\mathcal{T}_{1m}(z, w)$ is defined in the following way. For $\operatorname{Im} z > \frac{in}{4g}$ and $\operatorname{Im} w > \frac{im}{4g}$ we define it using the usual Beisert-Eden-Staudacher dressing factor σ [10]:

$$\mathcal{T}_{nm}(z,w) = \frac{2}{2\pi i} \frac{d}{dw} \log \sigma \left(x(z + \frac{in}{4g}), x(z - \frac{in}{4g}), x(w + \frac{im}{4g}), x(w - \frac{im}{4g}) \right) + \sum_{k=-\frac{n-1}{2}}^{\frac{n-1}{2}} \left(\mathcal{B}^{(1m)}(z + \frac{ik}{2g}, w) - \mathcal{R}^{(1m)}(z + \frac{ik}{2g}, w) \right) .$$
(A.7)

The function $\mathcal{T}_{1m}(z, w)$ has the branch-points at $z = \pm 1 + \frac{in}{4g}$ and $w = \pm 1 + \frac{im}{4g}$. One should analytically continue between them in z and w. Defined in this way function has four branch-cuts going to infinity in z variable starting at $z = \pm 1 \pm \frac{in}{4g}$ and four branch-cuts going to infinity in w variable starting at $w = \pm 1 \pm \frac{im}{4g}$ [52, 19]. In (2.57) we also use the notation $\mathcal{R}^{(10)\text{ph,mir}}$ and $\mathcal{T}_{1m}^{\text{ph,mir}}$, which means that one should

In (2.57) we also use the notation $\mathcal{R}^{(10)\text{ph,mir}}$ and $\mathcal{T}_{1m}^{\text{ph,mir}}$, which means that one should take $\mathcal{R}^{(10)}(z,w)$ (or $\mathcal{T}_{1m}(z,w)$) and then analytically continue it in the first argument along a path going around the branch point $z = 1 + \frac{i}{4g}$. For $\mathcal{R}^{(10)}(z,w)$ it simply results in the replacement $x(z \pm \frac{i}{4g}) \to x^{\text{ph}}(z \pm \frac{i}{4g})$.

In the main text we also use the following generalized Baxter polynomials

$$R^{(\pm)}(z) = \prod_{j=1}^{M} \frac{x(z) - x^{\rm ph}(z_j \mp \frac{i}{4g})}{\sqrt{x^{\rm ph}(z_j \mp \frac{i}{4g})}} , \quad B^{(\pm)}(z) = \prod_{j=1}^{M} \frac{1/x(z) - x^{\rm ph}(z_j \mp \frac{i}{4g})}{\sqrt{x^{\rm ph}(z_j \mp \frac{i}{4g})}} .$$
(A.8)

They are complex conjugates one of another $B^{\pm}(z) = \overline{R^{\mp}(z)}$.

B. Explicit expressions for *Y*'s

One can copy the expressions below directly to *Mathematica* from .pdf. Here we denote $d = \Delta$, $fb = \bar{f}$, $Ym[a_] = Y_{\bullet_a}$, $Yp[a_] = Y_{\triangle_a}$, $Yb[s_] = Y_{\bigcirc_s}$

```
sbf={e->f^2 d, T->f/fb,
S->((f-1)^2 fb(d fb-1)^2)/(f(fb-1)^2(d f-1)^2),
U->((f-1)^2(fb-1)^2)/(f^2 fb^2(d f-1)^2(d fb-1)^2),
A->1/(1-d f)+1/(1-d fb)+2/(d-1)+f/(f-1)+1/(fb-1),
F->-((2(f-1)^2(fb-1)^2)/((d-1)^3 f fb))-(f+fb-2 f fb)^2/((d-1)f fb)
-((f-1)(fb-1)((5 f-3)fb-3 f+1))/((d-1)^2 f fb)-f fb+2};
```

```
\begin{split} &Ym[a_] = (U \ S(T-1)^2 \ T \ e^a(e-T)^2(T^a(e-T^2)) (U \ S \ e^a(a+2)-T) - (e-1)U \ e^a(a+2) + \\ &(e-1)S \ T^a(2a) \ T^a)^2) / ((-(2 \ U \ T \ e^a(a+1)) (T-S \ T^a)(S \ T^a \ T-1)) \\ &+T^2(S \ T^a(2a)(T-1)^2 - U \ T \ e^a(S \ T^a-1)^2) + U \ e^a(a+2)(U \ S(T-1)^2 \ e^a \\ &-T(S \ T^a-1)^2)) (-(U \ T \ e^a(a+4)(S \ T^a \ T^2-1)^2) - U \ T^a \ e^a(a+2)(S \ T^a \ T^2-1)^2 \\ &+2 \ U \ T^2 \ e^a(a+3)(S \ T^a \ T-1)(S \ T^a \ T^3-1) + U^2 \ S(T-1)^2 \ e^a(2 \ a+6) \\ &+S \ T^a(2a)(T-1)^2 \ T^a)) / .sbf; \end{split}
```

```
\begin{split} & Yp[a_]=-(((e-1)(e-T^2) T^a(-(U T e^(a+3)(S T^(a+1)-1)^2)-U T^3 e^(a+1) \\ & (S T^(a+1)-1)^2+2 U T^2 e^(a+2)(S T^a-1)(S T^(a+2)-1)+U^2 S(T-1)^2 \\ & e^(2 a+4)+S(T-1)^2 T^(2 a+4)))/(((e-T^2) T^(a+1)(U S e^(a+3)-T)-(e-1) U \\ & e^(a+3) +(e-1) S T^(2 a+5))(U e^(a+2)(T-S T^a)+e T(S T^(a+1)-1) \\ & (U e^a-T^a)+T^(a+2)(S T^a-T))))/.sbf; \end{split}
```

 $Yb[s_]=(s-A)^2-1/.sbf;$

```
Y22=(d<sup>2</sup> f<sup>4</sup>(fb-1)<sup>2</sup>(d fb-1)<sup>2</sup>- d f<sup>3</sup>(d(d+1)fb<sup>2</sup>-4 d fb+d+1)(fb(d(2 fb-3)
-1)+2)+f<sup>2</sup>(d(d(d(d+4)+1)fb<sup>4</sup>-8 d(2 d+1)fb<sup>3</sup>+2(d(11-(d-7)d)+1)fb<sup>2</sup>
-8(2 d+1)fb+d+4)+1)-f(d(d+1)fb<sup>2</sup>-4 d fb+d+1)(fb(d(2 fb-3)-1)+2)+(fb-1)<sup>2</sup>
(d fb-1)<sup>2</sup>)/((d-1)f fb(d fb<sup>2</sup>((d-3)d f<sup>2</sup>-2(d-2)d f+d+2 f-3)-fb(d(2((d-2)
d-1)f<sup>2</sup>-3(d-1)d f+2 d+7 f-4)+f)+(d-3)d f<sup>2</sup>- 2(d-2)d f+d+2 f+2 fb-3));
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Y11=F/Y22/.sbf;

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