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Hagedorn Temperature of AdS$_5$/CFT$_4$ via Integrability

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We establish a framework for calculating the Hagedorn temperature of AdS$_5$/CFT$_4$ via integrability. Concretely, we derive the thermodynamic Bethe ansatz equations that yield the Hagedorn temperature of planar $\mathcal{N} = 4$ super Yang-Mills theory at any value of the ’t Hooft coupling. We solve these equations perturbatively at weak coupling via the associated $Y$ system, confirming the known results at tree level and one-loop order as well as deriving the previously unknown two-loop Hagedorn temperature. Finally, we comment on solving the equations at finite coupling.

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Introduction.—According to the AdS/CFT correspondence [1], $\mathcal{N} = 4$ super Yang-Mills (SYM) theory on $\mathbb{R} \times S^3$ is dual to type IIB string theory on AdS$_5 \times S^5$. This duality should in particular relate the phase transitions, critical behavior, and thermal physics of the theories.

One interesting example of a critical behavior is the Hagedorn temperature. In the planar limit of $\mathcal{N} = 4$ SYM theory on $\mathbb{R} \times S^3$, the origin of the Hagedorn temperature $T_H$ is the confinement of the color degrees of freedom due to the theory being on a three-sphere. This enables the theory to have a phase transition that bears resemblance to the confinement-deconfinement phase transition in QCD or pure Yang-Mills theory [2,3].

The Hagedorn temperature is the lowest temperature for which the planar partition function $Z(T)$ diverges. Via the state-operator correspondence, the partition function can be reexpressed in terms of the dilatation operator $D$ of $\mathcal{N} = 4$ SYM theory on $\mathbb{R}^4$:

$$Z(T) = tr_{R \times S^3}[e^{-H/T}] = tr_{R^4}[e^{-D/T}],$$

where we have set the radius of $S^3$ to 1. States correspond to gauge-invariant operators consisting of one or more trace factors. The energies correspond to the scaling dimensions of the operators, as measured by the dilatation operator. In the planar limit, the scaling dimensions of multitrace operators are entirely determined by those of their singletrace factors, and the latter can be enumerated via Pólya theory to determine the partition function and thus the Hagedorn temperature in the free theory [4]. This procedure was later generalized to one-loop order and to the case of nonzero chemical potentials [5–9].

On the string-theory side, the Hagedorn temperature occurs due to the exponential growth of string states with the energy present in tree-level string theory. For interacting string theory, it is connected to the Hawking-Page phase transition [10]. This suggests that the confinement-deconfinement transition on the gauge-theory side is mapped on the string-theory side to a transition from a gas of gravitons (closed strings) for low temperatures to a black hole for high temperatures. In particular, the Hagedorn temperature on the gauge-theory and string-theory sides of the AdS/CFT correspondence should also be connected [3,4].

On the string-theory side, the Hagedorn temperature has been computed in $p p$-wave limits [11–14]. In Ref. [15], the first quantitative interpolation of the Hagedorn temperature from the gauge-theory side to the string-theory side was made, exploiting a limit towards a critical point in the grand canonical ensemble [16]. This limit effectively reduces the gauge-theory side to the $\mathfrak{su}(2)$ sector with only the one-loop dilatation operator surviving, which enables one to match the Hagedorn temperature of the gauge-theory side to that of string theory on a $p p$-wave background via the continuum limit of the free energy of the Heisenberg spin chain.

A hitherto unrelated but very powerful property of planar $\mathcal{N} = 4$ SYM theory is integrability, see Refs. [17,18] for reviews. It amounts to the existence of an underlying two-dimensional exactly solvable model, which reduces to an integrable sigma model at strong coupling and to an integrable spin chain at weak coupling. Via integrability, the planar scaling dimensions of all singletrace operators can in principle be calculated at any value of the ’t Hooft coupling $\lambda = g^2_{\text{YM}} N$, allowing for a smooth interpolation between weak and strong coupling results. In practice, however, the calculation for each operator is so involved that summing the results for all operators to obtain the partition function $Z(T)$ seems prohibitive.

In this Letter, we show how to use integrability to compute the Hagedorn temperature at any value of the
We present them in the form of a Hagedorn temperature at any value of the $n$-dynamic Bethe ansatz (TBA) equations that determine the limit of the free energy of the spin chain associated with indeed, we find a direct relation between the continuum limit of the free energy of the spin chain associated with planar $\mathcal{N} = 4$ SYM theory and the Hagedorn temperature. Using the integrability of the model, we derive thermodynamic Bethe ansatz (TBA) equations that determine the Hagedorn temperature at any value of the ’t Hooft coupling. We present them in the form of a $Y$ system in Eqs. (11)–(23). As a first application, we solve them in the constant case as well as perturbatively at weak coupling, confirming the known tree-level and one-loop Hagedorn temperature. Moreover, we determine the previously unknown two-loop Hagedorn temperature:

$$T_H = \left( \frac{1}{2} \log(2 + \sqrt{3}) + \frac{1}{2} \log(2 + \sqrt{3}) \right) g^2 + \left( -\frac{86}{\sqrt{3}} + \frac{24 \log(12)}{\log(2 + \sqrt{3})} \right) g^4 + \mathcal{O}(g^6),$$

where $g^2 = \lambda/(16\pi^2)$.

**TBA equations for the Hagedorn temperature.**—In the following, we relate the Hagedorn temperature to the spin-chain free energy and derive TBA equations for the latter.

The Hagedorn temperature from the free energy of the spin chain: In the planar limit, the scaling dimensions of multitrace operators are completely determined by the scaling dimensions of their singletrace factors. The partition function $Z(T)$ is then entirely determined by the singletrace partition function $Z(T)$. Splitting the dilatation operator into a classical and an anomalous part as $D = D_0 + \delta D$, we can write

$$Z(T) = \sum_{n=2}^{\infty} e^{-\frac{2}{m} F_n(T)},$$

where

$$F_m(T) = -T \frac{2}{m} \log \left( \text{tr}_{\text{spin-chain,D}_0=\frac{m}{2}} \left[ e^{-\delta D(T)} \right] \right)$$

is the spin-chain free energy per unit classical scaling dimension for fixed $D_0 = (m/2)$. The multitrace partition function $Z(T)$ is then given by

$$Z(T) = \exp \sum_{n=1}^{\infty} \sum_{m=2}^{\infty} (-1)^{m(n+1)} e^{-\frac{2}{m} (n+F_n(T))/n},$$

where the alternating sign takes care of the correct statistics. The Hagedorn singularity is the first singularity of $Z(T)$ encountered raising the temperature from zero. It arises from the $n = 1$ contribution to the sum over $n$, i.e., from the infinite series

$$\sum_{m=2}^{\infty} e^{-\frac{2}{m} (1+F_n(T))},$$

where each term in the series is finite as $F_m(T)$ only includes a finite number of states. We can use Cauchy’s root test to assess when this series diverges. To this end, we compute the $m$th root of the absolute value of the $m$th term and take the large $m$ limit, giving

$$r = \lim_{m \to \infty} e^{-\frac{2}{m} (1+F_n(T))} = e^{-\frac{2}{m} (1+F(T))},$$

where

$$F(T) = \lim_{m \to \infty} F_m(T)$$

is the thermodynamic limit of the free energy. The root test states that the series is convergent for $r < 1$ and divergent for $r > 1$. Thus, the Hagedorn temperature is determined from $r = 1$ or, equivalently, from

$$F(T_H) = -1.$$
Note that TBA equations for the direct theory were also considered in Refs. [27,30] but in different thermodynamic limits.

Y system: The TBA equations can be rephrased in terms of a Y system consisting of the functions \( Y_{a,s}(u) \), where \((a, s) \in M = \{(a, s) \in \mathbb{N}_{\geq 0} \times \mathbb{N} | a = 1 \vee |s| \leq 2 \vee s = a = 2 \} \). With some exceptions, they satisfy the equations

\[
\log Y_{a,s} = \log \left( \frac{1 + Y_{a+1,s-1}(1 + Y_{a+1,s+1})}{(1 + Y_{a-1,s})(1 + Y_{a+1,s})} \right) \ast s, \tag{11}
\]

where \( \ast \) denotes the convolution with \( s(u) = (2 \cosh \pi u)^{-1} \) on \( \mathbb{R} \) and the (inverse) \( Y \) functions with shifted indices are assumed to be zero when the shifted indices are not in \( M \). The \( Y \) functions are analytic in the strip with \( \text{Im}(u) < \frac{1}{2}|a - |s|| \). For the purpose of this Letter, the chemical potentials are set to zero. Hence, the \( Y \) system is symmetric, \( Y_{a,s} = Y_{a,-s} \), with boundary conditions

\[
\lim_{a \to \infty} Y_{a+1,s} = 1, \quad \lim_{n \to \infty} Y_{1,n+1} = 1, \tag{12}
\]

for \( s = 0, \pm 1 \). The first of the aforementioned exceptions to the equations (11) then is

\[
\log Y_{1,0} = -\rho \ast s + 2 \log(1 + Y_{1,1}) \ast s - \log(1 + Y_{2,0}) \ast s, \tag{13}
\]

where we have defined \( \ast \) and \( \ast \) as the convolutions on \((-2g, 2g)\) and \( \mathbb{R} \setminus (-2g, 2g) \), respectively. Similarly, the convolution with \( Y_{1,1} \) and \( Y_{2,2} \) in Eq. (11) for \((a, s) = (2, 1), (1, 2)\) is also understood to be \( \ast \). The source term \( \rho(u) \) is defined as

\[
\rho = \frac{\rho_0}{T} + 2 \log(1 + Y_{1,1})(1 + Y_{2,2}) \ast H_0 + 2 \sum_{m=1}^{\infty} \log(1 + Y_{m+1,1}) \ast (H_m + H_{-m}) + \sum_{m=1}^{\infty} \log(1 + Y_{m,0}) \ast \Sigma^m, \tag{14}
\]

where

\[
H_m(u, v) = \frac{i}{2\pi} \partial_u \log \frac{x(u - i0) - \frac{g^2}{x(v + im)}}{x(u + i0) - \frac{g^2}{x(v + im)}}, \tag{15}
\]

\[
\rho_0(u) = \begin{cases} 0, & \text{for } |u| \geq 2g, \\ 2\sqrt{4g^2 - u^2}, & \text{for } |u| < 2g. \end{cases} \tag{16}
\]

and the kernel

\[
\Sigma^m(v, u) = i \frac{2}{2\pi} \partial_v \left( \frac{\log R^2(x(v + \frac{im}{2}), x(u + i0))}{R^2(x(v + \frac{im}{2}), x(u - i0))} \right. \\
+ \log \frac{R^2(x(v - \frac{im}{2}), x(u - i0))}{R^2(x(v - \frac{im}{2}), x(u + i0))} \right) \tag{17}
\]

is given in terms of the dressing factor \( [29] \)

\[
\sigma^2(u, v) = \frac{R^2(x^+(u), x^+(v))}{R^2(x^+(u), x^-(v))} \frac{R^2(x^-(u), x^+(v))}{R^2(x^-(u), x^-(v))} \tag{18}
\]

with \( x^+(u) = x(u \pm \frac{\pi}{2}) \). When applied to a function of two arguments such as \( H_m(v, u), \ast, \ast, \ast \) are moreover understood as integrals over the respective intervals. The other exceptions to the equations (11) are the nonlocal equations

\[
\log Y_{1,1}, Y_{2,2}(u) = \sum_{m=1}^{\infty} \log(1 + Y_{m,0}(v)) \ast \Theta_m(v, u) \tag{19}
\]

with

\[
\Theta_m(v, u) = \frac{i}{2\pi} \partial_v \log \frac{x(u) - \frac{g^2}{x(v + im)}}{x(u) - \frac{g^2}{x(v - im)}} \tag{20}
\]

and

\[
\log \frac{Y_{2,2}}{Y_{1,1}} = \sum_{m=1}^{\infty} a_m \ast \log \frac{(1 + Y_{m+1,1})}{(1 + Y_{1,m+1})^2(1 + Y_{m,0})} \tag{21}
\]

with \( a_m(u) = n/[2\pi(u^2 + n^2/4)] \).

The free energy per unit scaling dimension is given by

\[
F(T) = -T \sum_{n=1}^{\infty} \int_{-\infty}^{\infty} du \theta_n(u) \log (1 + Y_n,0(u)). \tag{22}
\]

where

\[
\theta_n(u) = \frac{i}{2\pi} \partial_n \log \frac{x(u + \frac{im}{2})}{x(u - \frac{im}{2})}. \tag{23}
\]

Thus, the TBA equations (11)–(23) determine the Hagedorn temperature at any value of the 't Hooft coupling via Eq. (9).

**Solving the TBA equations.**—Let us now solve the TBA equations in the form of the \( Y \) system.

Constant solution via \( T \) system: At large spectral parameter \( u \), the \( Y \) system approaches a constant value. This means we can find a constant \( Y \) system that solves Eq. (11) for all \((a, s) \in M \setminus \{(1, 1), (2, 2)\}\) as well as Eq. (21). Note that we cannot impose Eq. (19) as it relates
the behavior at finite and large \(u\). Thus, we find a one-
parameter family of solutions with parameter \(z\). This
solution is most easily expressed in terms of a \(T\) system
consisting of the functions \(T_{a,s}\) with \((a,s) \in \hat{M} = \{(a,s) \in \mathbb{Z}_{\geq 0} \times \mathbb{Z} | a < |s| \leq 2\}\) and \(T_{a,s} = 0\) for \((a,s) \not\in \hat{M}\). The \(Y\) functions are expressed in terms of the \(T\) functions as

\[
Y_{a,s} = \frac{T_{a,s+1}T_{a,s-1}}{T_{a+1,s}T_{a-1,s}}.
\] (24)

In the constant case, the equations (11) imply the following
\(T\) system (Hirota) equations for all \((a,s) \in \hat{M}:

\[
T_{a,s}^2 = T_{a+1,s}T_{a-1,s} + T_{a,s+1}T_{a,s-1}.
\] (25)

The latter are solved by

\[
T_{a,0} = \frac{(1 - z)^{2a} a + 2z a^2}{1 + z} + (12 z^2 - 1) a + 6z^2),
\]

\[
T_{a,\pm 1} = (\pm 1)^{a} \frac{(1 - z)^{2a} a + 3z a^2}{1 + z} (a^2 + 3az + 3z^2),
\]

\[
T_{a,\pm 2} = \frac{z^{-a}}{1 + z} (1 - z)^{2a},
\] (26)

for \(a \geq |s|\), and

\[
T_{0,s} = 1,
\]

\[
T_{1,s} = \frac{z^{-s}}{z^3} \left[ |s| + \frac{1 - 3z^2}{2z} \right] (1 - z^{|s|},
\]

\[
T_{2,s} = \frac{z^{-2s}}{1 + z} \left[ 1 - \frac{1 - z^2}{1 + z} \right],
\] (27)

for \(|s| \geq a\). This solution is a special case of the most
general, \(\mathfrak{psu}(2,2|4)\) character solution of Eq. (25) in
Ref. [31].

Solution at zero coupling: In the limit of zero coupling,
\(g^2 = 0\), the source term \(\rho(u)\) in Eq. (13) vanishes [32], such
that the functions \(Y_{a,s}\) are constant for all \(u\). Hence, the
nonlocal equation (19) implies \(Y_{1,1}Y_{2,2} = T_{1,0} = 1\). We
can use this to determine the parameter \(z\) in the constant
solution for the \(T\) system above and thereby find the \(Y\)
system at zero coupling. Imposing \(T_{1,0} = 1\) is equivalent to
\(z = \pm 1/\sqrt{3}\). The negative solution has to be discarded as it
leads to a negative Hagedorn temperature. Thus, we
conclude that to zeroth order \(z = 1/\sqrt{3}\). Using Eqs. (9) and
(22), we find the zeroth-order Hagedorn temperature

\[
T_{H}^{(0)} = \frac{1}{2 \log(2 + \sqrt{3})},
\] (28)

which is in perfect agreement with Ref. [4].

Perturbative solution: We can also solve the TBA
equations in a perturbative expansion at weak coupling,
expanding the \(Y\) functions as

\[
Y_{a,s}(u) = \hat{Y}_{a,s}^{(0)} \left( 1 + \sum_{\ell=1}^{\infty} g^{2 \ell} Y_{a,s}^{(\ell)}(u) \right).
\] (29)

At one-loop order, the solution takes the form

\[
Y_{a,s}^{(1)}(u) = \hat{Y}_{a,s}^{(1)} + \sum_{k=0}^{\infty} e_{a,s,k}^{(1)} \theta_{2k+a+s}(u),
\] (30)

where \(\hat{Y}_{a,s}\) as well as \(e_{a,s,k}^{(1)}\) are constants. This follows from the
expansions

\[
\epsilon_0 \hat{Y}(u) = 4\pi g^2 \hat{Y}(u) + 2\pi g^4 \hat{Y}''(u) + \mathcal{O}(g^6),
\]

\[
\hat{Y}(u) = \sum_{m=0}^{\infty} (-1)^m a_{1+2m}(u),
\]

\[
\theta_{u}(u) = a_{n}(u) + g^2 a_{n}''(u) + \mathcal{O}(g^4),
\]

\[
\Theta_{m}(v, u) = a_{m}(u - v) - a_{m}(v)
\]

\[
+ g^2 \left( \frac{2}{u} a_{m}''(v) - a_{m}''(v) \right) + \mathcal{O}(g^4),
\] (31)

in combination with the convolution identity \(a_{n} \ast a_{m} = \)
a_{n+m} and the structure of the TBA equations. Inserting
Eq. (30) into the expansion of the TBA equations, we can
solve for the coefficients \(e_{a,s,k}^{(1)}\). The remaining one-loop
parameter in the constant solution can be fixed from
\((Y_{1,1}Y_{2,2})^{(1)}(0) = 0\), which follows from Eq. (19) and
the last expansion in Eq. (31). We find for the one-loop
Hagedorn temperature

\[
T_{H}^{(1)} = \frac{1}{\log(2 + \sqrt{3})},
\] (32)

which perfectly agrees with the result of Ref. [5].

At two-loop order, \(\rho \ast \hat{Y}\) in Eq. (13) receives contributions
from the one-loop solution \(Y_{a,s}^{(1)}(u)\) from the second and
third term in Eq. (14). They can be calculated using

\[
(a_{n} \ast H_{m} \ast \hat{Y})(u) = g^2 \frac{4}{(n + |m|)^2} s(u) + \mathcal{O}(g^4).
\] (33)

Note that the dressing kernel in the fourth term of Eq. (14)
vanishes at this loop order. The two-loop solution takes the
form
In this Letter, we have used the fact (9) that the spin-chain free energy determines the Hagedorn temperature $T_H$ at which the partition function diverges. The spin-chain free energy should however also determine the partition function in the vicinity of $T_H$, which should allow us to extract, e.g., critical exponents.

The partition function and Hagedorn temperature have also been studied in integrable deformations of $\mathcal{N} = 4$ SYM theory up to one-loop order [38], where it was found that although $Z(T)$ is different, $T_H$ is unchanged. It would be interesting to see whether this statement continues to hold at higher loop orders. Similarly, one might apply our framework to the three-dimensional $\mathcal{N} = 6$ superconformal Chern-Simons theory, which is known to be integrable as well.

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Outlook.—In this Letter, we have derived integrability-based TBA equations (11)–(23) that determine the Hagedorn temperature of planar $\mathcal{N} = 4$ SYM theory at any value of the ’t Hooft coupling. As an application, we have solved these equations perturbatively up to two-loop order. Our TBA equation can also be solved numerically at finite coupling, as was briefly discussed here but will be detailed on in a future publication [34]. Thus, they open up the door for an exact interpolation from weak to strong coupling, which, with the exception of Ref. [15], would be the first time for the case of thermal physics. Potentially, this could allow us to develop a better understanding of the phase structure of gauge theories and their dual gravitational theories in general.

For the spectral problem, the TBA equations have been recast into the form of the quantum spectral curve [35], which allows one to generate precision data at weak coupling [36] as well as at finite coupling [37]. We will report on a similar reformulation of our equations in a future publication [34]. Moreover, one can study the case of nonzero chemical potentials. We have generalized our method to this case as well, and we have solved the zeroth-order TBA equations for the case with chemical potentials turned on but corresponding still to a symmetric $Y$ system. We will report on this in a future publication as well [34].
[19] The fact that finite-size corrections are suppressed in $L$ is manifest. The suppression in the classical scaling dimension can, for example, be seen for so-called twist operators. Up to corrections in the classical scaling dimension, these operators are dual to Wilson loops, for which no finite-size effects occur [20,21].