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From interacting particles to equilibrium statistical ensembles

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We argue that a particle language provides a conceptually simple framework for the description of anomalous equilibration in isolated quantum systems. We address this paradigm in the context of integrable models, which are those where particles scatter completely elastically and are stable against decay. In particular, we demonstrate that a complete description of equilibrium ensembles for interacting integrable models requires a formulation built from the mode occupation numbers of the underlying particle content, mirroring the case of noninteracting particles. This yields an intuitive physical interpretation of generalized Gibbs ensembles, and reconciles them with the microcanonical ensemble. We explain how previous attempts to identify an appropriate ensemble overlooked an essential piece of information, and provide explicit examples in the context of quantum quenches.

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The last decade has seen experimental techniques in optical lattices establish themselves as a fruitful playground for testing paradigms of quantum statistical physics [1–4]. A hallmark achievement has been the observation of anomalous equilibration in ultracold bosonic condensates [5–9] where the nonergodic character of quantum dynamics is attributed to the fact that they lie in the vicinity of an integrable point [10–24]. It became clear that a new framework would be required to properly describe such anomalous relaxation processes. In this work, we advocate a shift from conventional interpretations by offering a simple conceptual framework centered around the notion of particles. We argue that the consideration of anomalous equilibration in isolated quantum many-body systems naturally leads to the study of integrable models since they possess stable particle excitations. The techniques of integrability in turn allow for a complete description of generalized equilibrium states.

Recent efforts on the subject have predominantly revolved around the notions of prethermalization and generalized Gibbs ensembles (GGE). By invoking locality and entropy extremization, the GGE was introduced as the canonical statistical ensemble in which the Hamiltonian is supplemented with an extensive amount of additional conservation laws stemming from integrability. Initial studies focused on noninteracting particles, for which the conserved operators forming the GGE are the single-particle mode numbers [25–34]. Formulating the GGE for genuinely interacting models turned out to be more elusive, however, due to the less obvious structure of the local conservation laws [35–45].

In this work, we address the problem from a different viewpoint, and exploit the fact that, in contrast to generic (ergodic) dynamical systems, integrable systems exhibit stable collective excitations which can be identified as particles. While in noninteracting systems particles do not experience mutual collisions, the distinguishing property of interacting integrable models is that particles undergo a completely elastic scattering without particle production or decay [46,47]. The scattering is thus nondiffractive, meaning that momenta of the outgoing particles are simply the rearranged initial momenta. Here, we convey how anomalous thermalization in integrable systems reflects the existence of particles and use this insight to put forward a complete and universal description of local equilibrium states.

It turns out that our perspective demystifies and clarifies the physical picture of the unconventional equilibration seen in nonergodic interacting models, unifying its description with the case of noninteracting theories. Another advantage of adopting particle interpretation is that it eliminates the need of rather technical and subtle concepts introduced in the previous literature, e.g., recently advocated “weaker forms” of locality such as quasilocality [23,48] and semilocality [17], and makes the logic of the “truncated GGE” obsolete [33]. These auxiliary concepts, we argue, only undesirably obfuscate a clean physical picture. The very existence of an extensive hierarchy of higher (local and nonlocal) conservation laws is in fact a direct manifestation of the completely factorizable scattering of particles. It is the particles that are the inherently local objects of integrable theories [46].

In a closed quantum system, relaxation towards local equilibrium is typically formulated as an initial value problem, the so-called quantum quench, where a question of main interest is to characterize local equilibria which emerge as quasistationary states in the unitary relaxation process. This task can be naturally explained in the particle language, owing to the fact that integrability ensures that particles have infinite lifetimes and undergo completely elastic scattering. This observation implies that the information about the particle content at the initial time gets preserved for arbitrary long times. As we shall shortly discuss, this not only suffices to fully characterize the content of the stationary state, but it also proves necessary.

Prescribing an initial state can be viewed as exciting a macroscopic number of particles which subsequently participate in a scattering process. A local equilibrium state emerges dynamically after long times. However, the latter can only be meaningfully defined for a thermodynamically large system, which allows for a simplified description due to two mechanisms: dephasing and eigenstate thermalization hypothesis (ETH). Dephasing amounts to discarding the information contained in the dynamical phases from eigenstates at the same energy density [39,49], which averages out in the course of the relaxation process. ETH states that individual
microstates which are related by nonextensive modifications of quantum numbers are locally indistinguishable and give the same (generalized) free energy and expectation values of local observables [50]. No distinction between microcanonical and canonical ensembles is necessary as the former are understood as an unbiased collection of eigenstates which share the same set of distributions of quantum numbers, which for thermodynamically large systems exactly matches the canonical description, with the density operators of the particles providing a labeling of the microcanonical shells [50].

The task of finding and parametrizing the stationary state after a quench therefore amounts to detect the particle distributions from the initial state. There exist various strategies for performing this technical step. For instance, the one pursued in the seminal works on interacting quenches was to employ the quench action method [39,49] (see Appendix E) and compute the leading thermodynamic contribution of the overlap coefficients between the initial state and Bethe eigenstates [42–44]. While this has been successfully accomplished for certain special states [51], it seems to pose a formidable task for generic states. This can be overcome with the approach of Ref. [52] which derives exact operator representations for the quench action method [39,49] (see Appendix E) and compares it to the canonical description, with the density operators of the operators necessary to characterize general equilibrium states. The shift of perspective which we advocate in this work overcomes these difficulties.

As the entire spectrum of an integrable model is characterized in terms of stable particles, the number operators of these particles provide the most natural complete set for the local conserving $\hat{\Pi}$ in Eq. (1). Specifically, expressing the ensemble in terms of particle density operators $\hat{\rho}_j(v)$, it takes the form

$$\hat{\Pi} = Z^{-1} \exp \left[ -\sum_j \lambda_j \hat{Q}_j \right],$$

where the index $j$ runs over distinct particle types, $v$ is a rapidity variable which parametrizes their momenta $p_j(v)$, and $\mu_j(v)$ are rapidity-dependent chemical potentials. The eigenvalue densities $\rho_j$ of the operators $\hat{\rho}_j$ completely determine the macrostates of the system [50]. Individual microstates drawn from the same macrostate differ only by rearrangements of the occupation of the $Lp_j(v)dv$ quantum numbers in a rapidity window $dv$, and thus have a negligible effect on the local correlations in the thermodynamic limit. The entropy takes the universal form given explicitly below [cf. Eq. (12)]. This is understood as the integrable analog of ETH used to characterize the thermalization of eigenstates in ergodic interacting systems [55–58].

**Heisenberg model.** We focus on the prototypical integrable model, the Heisenberg spin-$1/2$ chain,

$$\hat{H} = -\frac{1}{2} \sum_j [\tilde{S}_j \cdot \tilde{S}_{j+1}].$$

The particle content (cf. Appendix A), with respect to a fully polarized state, consists of magnons and $j$-particle bound states thereof, with momenta [59]

$$p_j(v) = -i \log \left[ (v + j \frac{i}{2})/(v - j \frac{i}{2}) \right].$$

Recently, a set of local charges $\hat{X}_j(v)$ for the Heisenberg model has been identified [60], which allows macrostates to be uniquely determined from the values of charge densities $X_j$ via string-charge duality [52]. This takes the form of a discrete wave equation

$$\rho_j = \Box X_j,$$

where $\Box$ is a d’Alembertian on $X_j(v)$ defined for a set of functions $f_j(v)$ as

$$\Box f_j = f_j^+ + f_j^- - f_{j-1} - f_{j+1},$$

where $f_j^\pm(v) = f_j(v \pm i/2 \mp 0^+)$ and $0^+$ denotes an infinitesimal which acts as a regulator.

Although previous attempts at identifying a complete ensemble were built directly from a discrete basis of charges obtained from $\hat{X}_j$, the strategy they used was to incorporate the knowledge of charge densities to enforce constraints on the space of macrostates, thus tacitly bypassing the canonical form of Eq. (1) which is merely a formal series. Since it remains unclear how to unambiguously interpret it, we instead employ a continuous set of local charges $\hat{X}_j(v)$ and cast the most
The functions $\tilde{\rho}_j$ represent the densities of unoccupied modes, which obey the Bethe-Yang equations \[ \rho_j + \tilde{\rho}_j = a_j - a_{j,k} \rho_k, \] where the integral kernels $a_{i,k}$ are given by $a_{j,k} = G_{j,k} + G_{j,k+1}$. The dominant contribution to $Z$ in the thermodynamic limit is given by the saddle point, yielding a set of coupled nonlinear integral equations known as the TBA equations \[ \log Y_j = \mu_j + a_{j,k} \log \left( 1 + Y_k^{-1} \right), \] written in terms of $Y$ functions $Y_j = \tilde{\rho}_j / \rho_j$. Together, Eqs. (13) and (14) completely determine the equilibrium state. They provide the direct relationship between the set of chemical potentials $\mu_j$ and the set of macrostates $\rho_j$, demonstrating the completeness of the ensemble given in Eq. (2).

For practical purposes, it is useful to switch to an alternative formulation of the TBA equations. The infinite sum in Eq. (14) can be eliminated with the use of kernel identities, yielding a local form \[ \log Y_j = d_j + s \log \left( 1 + Y_{j-1} \right) \left( 1 + Y_{j+1} \right), \] where the source terms are \[ d_j = \mu_j - s \left( \mu_{j-1} + \mu_{j+1} \right), \] and the convolution kernel \[ s(v) = \left[ 2 \cosh(\pi v) \right]^{-1}. \] These are equivalent to Eq. (14) when supplemented with the large-$v$ asymptotics of the $Y$ functions. To proceed, care must be taken when inverting the convolution with $s$, as the pseudoinverse defined through \[ s^{-1} * f = f^+ + f^- \] has a nontrivial null space (see, e.g., \[ 66,67\]). In particular, while $s^{-1} * (s * f) = f$ for any function $f$, in general $s * (s^{-1} * f) \neq f$. We therefore decompose the source terms as \[ d_j = s * \lambda_j + d_j^\theta, \] where \[ \lambda_j = s^{-1} * d_j, \] and the “singular components” $d_j^\theta$ are annihilated by $s^{-1}$. Applying $s^{-1}$ to Eq. (15) and exponentiating, yields the modified $Y$-system relations \[ Y_j^+ Y_j^- = e^{\xi^1 (1 + Y_{j-1}) (1 + Y_{j+1})}. \] In the process, the singular components are not lost, but are instead encoded in the analytic data of the $Y$ functions: their poles and zeros in the physical strip $\mathcal{P}$, Equation (15) is reobtained by convolving the logarithm of Eq. (21) with $s$, after multiplying out the singularities $\xi$ with the functions \[ t(v;\xi) = \tanh \left( \frac{\xi v}{2} \right), \quad |\text{Im}(\xi)| < \frac{1}{2}, \] which satisfy $t^t t^- = 1$. As a result, the singular components are of the form \[ d_j^\theta(v) = \sum_a \log t(v;\xi_a^2) - \sum_b \log t(v;\xi_b^4), \] where the combinatorial weight is given by the Yang-Yang entropy density per mode [50].
where \( \{\xi^q_{\mu}\} \) and \( \{\xi^r_{\nu}\} \) are, respectively, the sets of zeros and poles of the \( Y \) function \( Y_j \) in \( \mathcal{P} \).

We now address the question of the generality of the charge-based GGE of Eq. (7). Using Eq. (10), this ensemble takes the form of Eq. (2) with

\[
\mu_j = G_{j,k} \cdot \lambda_k + \lambda_j = 0.
\]

Naively, the two ensembles could then be assumed equivalent. The crucial point, however, is that the Lagrange multipliers \( \mu_j \) are blind to the singular components \( d^s_j \). In fact, as \( \square \mu_j = \lambda_j \), the Lagrange multipliers precisely match the functions \( \lambda_j \) introduced in the decomposition of Eq. (19).

Put another way, the d’Alembertian \( \square \) inherits a null space from \( s^{-1} \), and \( \square \mu_j = \lambda_j \) demonstrates that the Lagrange multipliers \( \lambda_j \) of Eq. (7) cannot encode the components of \( \mu_j \) in this null space. In the following section, we provide evidence that singular component contributions are a generic feature of equilibrium states. To conclude, it is worthwhile to stress that the null space of \( \square \) does not cause a problem for the string-charge duality relations [Eqs. (5) and (10)] as the densities \( \rho_j \) on which \( \square \) acts already provide the full description of a macrostate.

### III. Examples

We now explicitly demonstrate this limitation of the charge-based GGE of Eq. (7) with some examples in the context of a quantum quench. Starting from an initial state \( |\Psi\rangle \), the equilibrium macrostate \( \rho_j \), reached in the long-time limit is determined from the expectation values of the particle density operators

\[
\rho_j(v) = \lim_{L \to \infty} L^{-1} \langle \Psi | \hat{\rho}_j(v) | \Psi \rangle.
\]

As described above, it is convenient to encode a macrostate in the \( Y \) functions, \( Y_j = \hat{\rho}_j / \rho_j \), the condition that the charge-based GGE is a valid ensemble is the absence of poles and zeros of these functions in the physical strip \( \mathcal{P} \).

The best studied quenches are those from the “dimer” state

\[
|D\rangle = \frac{1}{2^{L/2}} (|\uparrow \downarrow \rangle - |\downarrow \uparrow \rangle)^{\otimes L/2},
\]

and the Néel state

\[
|N\rangle = |\uparrow \downarrow \rangle^{\otimes L/2},
\]

for which the \( Y \) functions are explicitly known [70,71]. For the dimer quench these are [52]

\[
Y_j(v) = \frac{2j(j + 2)v^2}{[v + (j + 1)\frac{1}{2}] [v - (j + 1)\frac{1}{2}]}. \tag{28}
\]

The double zero at the origin immediately indicates that the Lagrange multipliers \( \lambda_j \) in the charge-based GGE are not sufficient to capture it. Indeed, the source terms

\[
d_j(v) = \log \tanh \left( \frac{\sqrt{v}}{2} \right)^2, \tag{29}
\]

determined through Eq. (15), are in the null space of \( s^{-1} \), from which it follows that \( \lambda_j = 0! \) We stress that the state is by no means close to the infinite-temperature Gibbs state in terms of local correlation functions. Previous works mistakenly assigned nontrivial \( \lambda_j \) for this state [48,70,71], resulting in the incorrect interpretation of the GGE. The important point is that due to the nontrivial null space of \( s^{-1} \), it is not permissible to define \( \lambda_j \) via Fourier transform, namely, \( \mathcal{F}[\lambda_j] \neq \mathcal{F}[d_j]/\mathcal{F}[s] \).

The Néel state is similar. Again, the source terms

\[
d_j(v) = (-1)^{j+1} \log \tanh \left( \frac{\sqrt{v}}{2} \right)^2 \tag{30}
\]

consist solely of singular components, indicating \( \lambda_j = 0 \) for this state also. Indeed, these two examples are members of a particular class of initial states, whose equilibrium states can be cast as the partition function of a vertex model for which the initial state provides an integrable boundary [72]. Generic initial states do not admit such a description.

We thus also consider more general product states. A complete analytic treatment is now out of scope, as the complexity of \( \rho_j \), computed via Eq. (25), grows quickly with \( j \). Nevertheless, the lowest \( Y \) functions are straightforwardly determined. A particular example is the state composed of alternating two-site domain walls \( |\uparrow \uparrow \downarrow \downarrow \rangle^{\otimes L/4} \). Explicit expressions are unwieldy (see Appendix D). They indicate nontrivial \( \lambda_j \), and that \( Y_1 \) has four zeros at \( \pm(0.382 \pm 0.234 i) \) and two poles at \( \pm0.155 i \). These analytic data indicate that again the resulting equilibrium state is not captured by the GGE of Eq. (7).

### IV. Conclusion

We have shown that a complete description of equilibrium states in interacting integrable models requires the ensemble to be constructed from the number operators of the model’s particle content. Any attempt to reduce the description inevitably leads to a loss of information, rendering a part of the equilibrium manifold inaccessible. The proposal naturally extends the established framework for noninteracting particles to the interacting integrable regime. The effect of nontrivial scattering in the interacting case is that excitations about an equilibrium state get nontrivially dressed by the correlations to the interacting integrable regime. The effect of nontrivial scattering in the interacting case is that excitations about an equilibrium state get nontrivially dressed by the correlations of the state (see, e.g., [50,73]).

Our hope is that the particle-based perspective which we advocate here creates a platform for the study of equilibration as one leaves the integrable points. The resulting loss of factorized scattering induces time scales up to which the particle-based description of ensembles may be expected to accurately capture statistical properties of quasistationary states, the so-called prethermalized regime [74-77].

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APPENDIX A: STRING HYPOTHESIS AND BETHE-YANG EQUATIONS

Any eigenstate of the isotropic Heisenberg spin chain of finite length $L$ is associated a unique set of $M$ rapidities. These are obtained from solutions to Bethe quantization conditions

$$e^{ip(v)L} \prod_{k=1}^{M} S_{1,k}(v - v_k) = -1, \quad v = \{v_j\}_{j=1}^{M}, \quad (A1)$$

where $M/L$ is the magnetization density,

$$p(v) = -i \log \left( \frac{u + i \frac{j}{2}}{u - i \frac{j}{2}} \right), \quad (A2)$$

and the scattering amplitude $S_{1,1}$ is a member of a complete set of scattering amplitudes

$$S_{j}(v) = \frac{v - j \frac{1}{2}}{v + j \frac{1}{2}}, \quad (A3)$$

$$S_{j,k} = S_{j-k,k} \prod_{m=0}^{\min(j,k)-1} S_{j-k+2m,m}. \quad (A4)$$

In the thermodynamic limit ($L \to \infty$, $M \to \infty$, keeping $M/L$ fixed), the solutions (rapidities) $v_j$, $j = 1, 2, \ldots, M$ (referred to as the Bethe roots) in the complex plane organize in a special way and permit to partition the entire spectrum in terms of particles. A central feature of generic integrable models is the formation of bound states, representing stable particles which exhibit elastic scattering without particle production. In terms of solutions $v_j$ ($j = 1, \ldots, M$) to Eq. (A1), these comprise of rapidities with nonzero imaginary parts which share common real parts. According to the string hypothesis [59,61,62], the spectra of a thermodynamically large system can be partitioned in terms of complex-valued solutions, which physically represent multimagnonic excitations called the strings. A $k$-string solution ($k \in \mathbb{N}$) centered at $v^*_k \in \mathbb{R}$ is parametrized as

$$\left\{ v^{k,l}_a \right\} = \left\{ v^*_a + (k + 1 - 2l) \frac{i}{2} : l = 1, 2, \ldots k \right\}. \quad (A5)$$

suppressing exponentially small deviations in system size $L$.

In the $L \to \infty$, the string centers become dense on real axis and condense. This allows to introduce distributions of Bethe $j$ strings $\rho_j$ and their hole counterparts $\bar{\rho}_j$ (unoccupied solutions allowed by Bethe quantization condition). In the thermodynamic limit, the quantization condition (A1) gets replaced by linear integral equations which are customary called the Bethe-Yang equations [50,59,61]

$$\rho_j + \bar{\rho}_j = a_j - a_{j,k} \ast \rho_k, \quad (A6)$$

where here and below we assume summation convention over repeated indices, and introduce convolution operation $(f \ast g)(v) = \int_0^L df(t) f(v - t)g(t)$. The integral kernels $a_j$ and $a_{j,k}$ represent the “kinematic data”, i.e., encode full scattering data among all distinct types of strings. Specifically, they are the derivatives of scattering phase shifts

$$a_j(v) = \frac{1}{2\pi i} \frac{\partial}{\partial v} \log S_j(v), \quad (A7)$$

$$a_{j,k}(v) = \frac{1}{2\pi i} \frac{\partial}{\partial v} \log S_{j,k}(v). \quad (A8)$$

APPENDIX B: LOCAL CHARGES AND THERMODYNAMIC SPECTRA

In accordance with the standard practice in Yang-Baxter integrable models [53,54], we introduce an infinite set of quantum transfer operators

$$\hat{T}_j(v) = Tr_{V_j} \hat{L}_j^{(1)}(v) \hat{L}_j^{(2)}(v) \ldots \hat{L}_j^{(L)}(v), \quad (B1)$$

defined as traces over $L$-fold spatially ordered products of Lax operators $\hat{L}_j^{(l)}(u)$ acting in the tensor product $\mathcal{H} \otimes V_j$, where $\mathcal{H} \cong V^{\otimes L}$ is the Hilbert space of the spin chain. Specifically, Lax operators take the form

$$\hat{L}_j^{(l)}(v) = 1 \otimes 1 \otimes \ldots \otimes \hat{L}_j(v) \otimes 1 \otimes \ldots \otimes 1, \quad (B2)$$

$$\hat{L}_j(v) = v 1 \otimes 1 + i \sum_{\alpha = \{x,y,z\}} \hat{S}_\alpha \otimes \hat{S}_j^\alpha, \quad (B3)$$

where $\hat{S}_\alpha^\alpha$ denotes spin-$j/2$ operators acting in $V_j$ which enclose the algebraic relations $[\hat{S}_j^\alpha, \hat{S}_\beta^\beta] = i \epsilon_{\alpha\beta\gamma} \hat{S}_j^{\gamma}$. Conserved operators $\hat{T}_j$ mutually commute,

$$[\hat{T}_j(v), \hat{T}_j(v')] = 0, \quad (B4)$$

for all values of $j,j' \in \mathbb{N}$ and $v' \in \mathbb{C}$, as implied by the Yang-Baxter relation. A set of local charges is then defined through their logarithmic derivatives [48,60]

$$\hat{X}_j(v) = \frac{1}{2\pi i} \frac{\partial}{\partial v} \log \hat{T}_j^{(1)}(v), \quad (B5)$$

where $\phi = T_0^{(+)}$ and the spectral parameter $v$ should be restricted to the physical strip $\mathcal{P}$, defined as the strip

$$\mathcal{P} = \{ v \in \mathbb{C} : |\text{Im}(v)| < \frac{1}{2} \}. \quad (B6)$$

We note that the conserved operators $\hat{X}_j(v)$ are only defined on $\mathcal{P}$ [23], and become singular as the boundary is approached. The Heisenberg spin-chain Hamiltonian $H$ is proportional to the charge $\hat{X}_1(0)$. By convention, we adopt $\hat{H} \equiv \hat{X}_1(0)$.

The action of $\hat{X}_j$ on Bethe eigenstates in the $L \to \infty$ limit has been obtained in Refs. [48,52]. Expressed in terms of Bethe root densities $\rho_j$ they take the compact form

$$\hat{X}_j = G_{j,k} \ast \rho_k, \quad (B7)$$

with

$$G_{j,k} = \sum_{m=1}^{\min(j,k)} a_{j-k+2m,m}. \quad (B8)$$

Inverting relation (B7) amounts to find an operator $\square$ such that $\square G_{j,k}(v) = \delta_{j,k} \delta(v)$. To this end, we employ kernel identities [59]

$$a_j - I_{j,k} \ast a_k = \delta_{j,k}, \quad (B9)$$

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where the convolution kernel reads as
\[
s(v) = \frac{1}{2 \cosh(\pi v)},
\]
and the incidence matrix \( I_{j,k} = \delta_{j,k-1} + \delta_{j,k+1} \) expresses how distinct types of particles interact with each other. As Eq. (B9), special attention has to be paid to the pole structure of \( a_1 \). To properly account for the singularity at the boundary of \( \mathcal{P} \) we introduce a pseudoinverse
\[
s^{-1} \ast f = f^+ + f^-.
\]

Using the pseudoinverse we define a discrete d’Alembertian
\[
\square y = y^+ - y^-.
\]

This means that \( s^{-1} \) is only a left inverse of \( s \),
\[
s^{-1} \ast (s \ast f) = f,
\]
while in general there exist functions \( f \) for which
\[
s \ast (s^{-1} \ast f) \neq f.
\]

A subtle point here is that any constant term in \( \mu_j \), does not enter in \( d_j \) as 1 \( \ast s = \frac{1}{2} \). Equation (C3) can be reobtained by casting Eq. (C5) as
\[
\log Y_j = d_j + I_{j,k} \mu_k \ast log (1 + Y_k),
\]
referred to as the local TBA equations. The source terms \( d_j \) are related to the chemical potentials
\[
d_j = \mu_j - I_{j,k} s \ast \mu_k.
\]

A naive “integration” of the (modified) TBA equations by virtue of Cauchy theorem is not possible. A standard approach to perform thermodynamic considerations is to work under the string hypothesis and introduce a discrete set of \( Y \) functions,
\[
Y_j = \frac{\tilde{\rho}_j}{\rho_j},
\]
given as ratios of hole and particle densities for each string species. In the thermodynamic limit (by accounting for the Yang-Yang entropy as described in the text), a functional representation for the partition sum yields an infinite set of coupled nonlinear integral equations
\[
\log Y_j = \mu_j + a_{j,k} \ast \log (1 + Y_k^{-1}).
\]

### APPENDIX C: CANONICAL AND LOCAL FORM OF TBA EQUATIONS

We now turn our attention to the statistical ensemble
\[
\hat{\rho} = \mathcal{Z}^{-1} \exp \left[ -\sum_j \int d\nu \mu_j(v)\hat{\rho}_j(v) \right],
\]
where the normalization \( \mathcal{Z} = \text{Tr} \hat{\rho} \) represents the partition sum. A standard approach to perform thermodynamic considerations is to work under the string hypothesis and introduce
\[
\lim_{v \to -\infty} \tilde{Y}_j(v) = 1. 
\]
Now convolving with \( s \) after taking the logarithm gives back the local TBA equations
\[
\log Y_j = d_j^0 + s \star \lambda_j + I_{j \perp} \star \log(1 + Y_k), \tag{C11}
\]
where we introduced the splitting of the source terms \( d_j = s \star \lambda_j + d_j^0 \), with the singular part
\[
\tilde{d}_j^0 = \sum_a \log \left( \frac{e^{z_a/\beta}}{s^a \beta} \right) - \sum_b \log \left( \frac{e^{z_b/\beta}}{s^b \beta} \right). \tag{C12}
\]
It is worthwhile remarking that the physical solutions correspond to real valued \( Y_j(v) \) on the real axis, requiring the singularities \( \xi \) to always appear in complex-conjugate pairs.

Going in the opposite direction, i.e., transforming Eq. (C11) is straightforwardly obtained by applying \( s^{-1} \) to both sides and subsequently taking the exponent. In the process, the information from the singular part \( d_j^0 \) gets transferred into the analytic structure of \( Y \) functions, while \( \lambda_j \) enter the modified \( Y \)-system (C8) as the nonuniversal (i.e., node-dependent) part. Notice that setting \( \lambda_j \equiv 0 \) recovers the universal form which is common in the literature, governing states whose physical input (TBA source terms) consist solely from the null-space components.

\textit{Gibbs equilibrium.} The cleanest example of a state which fulfills the modified set of functional relations (C8) is the canonical Gibbs equilibrium, corresponding to bare dispersions \( \mu_j^{\text{Gibbs}} = \beta a_j \). Due to absence of zero modes, this is equivalent to have \( \lambda_j^{\text{Gibbs}} = \mu_j^{\text{Gibbs}} \), with
\[
\lambda_j(v) = \beta \delta_{j,1}, \tag{C13}
\]
restoring the well-known TBA source term
\[
\tilde{d}_j^{\text{Gibbs}} = s \star \lambda_j^{\text{Gibbs}} = \beta s. \tag{C14}
\]

\section*{Appendix D: Exact Solutions of Quantum Quenches}
Below we explain how to detect zero modes given a macrostate \( \rho_j^\psi \). In the context of a quench protocol \( \rho_j^\psi \) can be understood as equilibrium states of a quench with the initial condition \(|\Psi\rangle\). We first consider two simple product states which have represented toy examples in previous studies of quantum quenches in the Heisenberg model \([41,43,44,48,52,70–72]\).

\textit{Dimer state.} We first consider the so-called dimer state
\[
[D] = \frac{1}{2^{L/2}} (|\uparrow \downarrow\rangle - |\downarrow \uparrow\rangle)^{\otimes L/2}. \tag{D1}
\]
The associated macrostate \( \rho_j \) expressed in terms of \( Y \) functions reads as \([52]\)
\[
Y_j(v) = \frac{4v^2(12v^2 + 19)}{(v + \frac{1}{2})(v - \frac{1}{2})(v + \frac{1}{2})(v - \frac{1}{2}) - 16v^2(v - \frac{1}{2})(v + \frac{1}{2})}. \tag{D2}
\]
These can be quickly verified to obey the standard form of \( Y \)-system relations \([65,78,79]\)
\[
Y_J Y_J = (1 + Y_{j-1})(1 + Y_{j+1}), \tag{D3}
\]
representing a specialization of Eq. (C8) when \( \lambda_j \equiv 0 \). These relations have a purely group-theoretic origin and represent the kinematic input to the problem. The physical input of the solution is encoded in analytic properties of \( Y_j \) in the complex plane. In the present case it consists of a double zero at the origin \( v = 0 \) and a pair of simple poles at \( \pm (j + 1)\frac{1}{2} \), while the large-\( v \) asymptotics matches that of the infinite-temperature Gibbs equilibrium, that is, \( Y_j(v) \sim j(j + 2) \) as \( |v| \to \infty \). The only physically relevant input lies inside the strip \( \mathcal{P} \), where we find for every function \( Y \) a double zero at the origin, yielding the local TBA source terms
\[
\tilde{d}_j(v) = \log \left( \frac{v}{v + \frac{1}{2}} \right)^2. \tag{D4}
\]
The TBA source terms obey the fusion property
\[
\square g_j = 0, \tag{D5}
\]
where
\[
g_j(v) = \log \left( \frac{4v^2 \left( v + \frac{1}{2} \right) \left( v - \frac{1}{2} \right)}{\left( v + \frac{1}{2} \right) \left( v - \frac{1}{2} \right)} \right). \tag{D6}
\]
\textit{Néel state.} The antiferromagnetic (Néel) state
\[
[N] = |\uparrow \downarrow\rangle^{\otimes L/2} \tag{D7}
\]
represents another particularly simple example which can be treated exactly (see, e.g., \([43,70]\)). The initial two \( Y \) functions read as explicitly \([52]\)
\[
Y_1(v) = \frac{4v^2(12v^2 + 19)}{(v + \frac{1}{2})(v - \frac{1}{2})(v + \frac{1}{2})(v - \frac{1}{2}) - 16v^2(v - \frac{1}{2})(v + \frac{1}{2})}, \tag{D8}
\]
\[
Y_2(v) = \frac{v^2(v + \frac{1}{2})(v - \frac{1}{2})}{(v + \frac{1}{2})(v - \frac{1}{2})(v + \frac{1}{2})(v - \frac{1}{2})}. \tag{D9}
\]
whereas the higher ones can be computed from the \( Y \)-system relations (D3). Importantly, the analytic data of \( Y_1 \) inside \( \mathcal{P} \) are now given by double zeros (poles) for index \( j \) being odd (even). This implies
\[
\mu_1(v) = \log \left( \frac{16v^2}{(v + \frac{1}{2})(v - \frac{1}{2})} \right). \tag{D10}
\]
while the higher source terms follow from the fusion condition
\[
\mu_j = \sum_{k=1}^j \mu_1^{j+1-2k}. \tag{D11}
\]

It is worth stressing that the two examples given here in fact represent atypical initial conditions. This is attributed to the fact that the TBA source have only components from \( d_j^0 \), implying \( \lambda_j \equiv 0 \). General initial states \(|\Psi\rangle\) on the other hand involve nonvanishing \( \lambda_j \) or, put differently, are given by functions \( \mu_j \) which do violate the fusion condition (D11). A practical drawback of this is that a full specification of a macrostate \( \rho_j \) goes beyond the knowledge of only \( \mu_1 \).

\textit{Generic states.} Here, we analyze a nontrivial periodic state of 2-spin ferromagnetic domain walls
\[
|\Phi\rangle = |\uparrow \downarrow \downarrow \uparrow\rangle^{\otimes L/4}, \tag{D12}
\]
representing one of the simplest generic initial states. This state has already appeared in previous studies of quantum
Expressions for the first few $Y$ functions can be computed analytically, but are already rather formidable. The initial $Y$ function for instance reads as

$$Y_1(v) = \frac{(64v^6 + 112v^4 - 20v^2 + 5)(192v^6 + 784v^4 + 1124v^2 + 491)}{(4v^2 + 1)(4v^2 + 5)(256v^8 + 1024v^6 + 1696v^4 + 1248v^2 + 29)}$$  \hfill (D13)

with four zeros at $\pm (0.382 \pm 0.234i)$ and two poles at $\pm 0.155i$, in $P$. It can furthermore be easily verified analytically that the universal form of the $Y$-system (D3) is now no longer satisfied (see also [80]). This specifically means that after subtracting the singular components $d_s^j \lambda_j$ from the stationary part of the quench action systems, the steady-state limit of the evolution is entirely observables in thermodynamically large integrable many-body systems. The steady-state limit of the evolution is entirely determined from the stationary part of the quench action functional

$$Z_{QA} = \int \mathcal{D}[\rho_j] \exp(-L S_{QA}).$$  \hfill (E1)

APPENDIX E: QUENCH ACTION

The quench action approach [39,49] has been introduced as a functional integral approach towards time evolution of local observables in thermodynamically large integrable many-body systems. The steady-state limit of the evolution is entirely determined from the stationary part of the quench action functional

$$Z_{QA} = \int \mathcal{D}[\rho] \exp(-L S_{QA}).$$

$$S_{QA} = \sum_j \int_{\mathbb{R}} dv \left(\mu_j(v) \rho_j(v) + s_j(v)\right).$$  \hfill (E2)

The state-dependent terms $\mu_j$ represent the mode decomposition of the density of the logarithmic overlap between $|\Psi\rangle$ and a general macrostate, namely,

$$\sum_j \int_{\mathbb{R}} dv \mu_j(v) \rho_j(v) = \lim_{L \to \infty} L^{-1} \log |\langle \Psi | [\nu_j] \rangle|^2.$$  \hfill (E3)

Saddle-point evaluation of the functional $Z_{QA}$ yields a sought-for macrostate $\rho_j$. As the choice of notation already suggests, $\mu_j$ of Eq. (E3) are nothing but the chemical potentials corresponding to individual particle modes. This can be most easily be seen by comparing Eq. (E3) with the form of Eq. (C1), implying that $Z_{QA}$ coincides with the partition sum $\mathcal{Z}$ of the generalized Gibbs ensemble $\mathcal{G}$.