Letter

Resonance triplet dynamics in the quenched unitary Bose gas

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The quenched unitary Bose gas is a paradigmatic example of a strongly interacting out-of-equilibrium quantum system, whose dynamics become difficult to describe theoretically due to the growth of non-Gaussian quantum correlations. We develop a conserving many-body theory capable of capturing these effects, allowing us to model the postquench dynamics in the previously inaccessible time regime where the gas departs from the universal prethermal stage. Our results show that this departure is driven by the growth of strong lossless three-body correlations, rather than atomic losses, thus framing the heating of the gas in this regime as a fully coherent phenomenon. We uncover the specific few-body scattering processes that affect this heating and show that the expected connection between the two-body and three-body contacts and the tail of the momentum distribution is obscured following the prethermal stage, explaining the absence of this connection in experiments. Our general framework, which reframes the dynamics of unitary quantum systems in terms of explicit connections to microscopic physics, can be broadly applied to any quantum system containing strong few-body correlations.

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Introduction. It is a well-known fact that the quantum many-body problem scales exponentially with the number of particles. Hence, despite the rapid rise in available computing power over the last few decades, exact descriptions typically remain out of computational reach [1,2]. In many cases, however, simplifying assumptions can be made about the underlying processes of the system, namely, that they obey Gaussian statistics and can be described by long-lived quasiparticles and collective motions of the medium using just a few degrees of freedom [3]. Such simplifications are at the heart of our understanding of a wide range of quantum phenomena, ranging from Fermi liquids and polarons to superfluids and conventional superconductors [4–10].

In realistic quantum systems, however, quasiparticles attain a finite lifetime due to the presence of non-Gaussian correlations, which allow collisions and decay into the many-body continuum [11–13]. In extreme cases, these correlations lead to such short lifetimes that the quasiparticle spectrum is no longer well defined and the standard simplifications do not apply. The theoretical challenge of understanding such strongly correlated systems is important in many branches of physics, such as condensed matter [14–20], ultracold gases [21–29], nuclear physics [30–32], and quantum technologies [33–38].

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Due to the large degree of control over the interaction strength, ultracold atomic gases provide a versatile quantum simulator for probing non-Gaussian physics [39–41]. Additionally, the underlying microscopic physics can be encoded exactly in many-body models of these systems, enabling quantitative comparisons between theory and experiment [42,43]. At the frontier of such approaches is the description of ultracold atomic many-body systems featuring nonperturbative few-body effects [44–49]. Here, a series of experiments exploring the quench of a degenerate Bose gas to the unitary regime [50–53], where interactions are as strong as allowed by quantum mechanics, raise important fundamental questions regarding the interplay of integrability, ergodicity, and few-body correlations in strongly correlated systems far out of equilibrium. Immediately following the quench, the dynamics are Gaussian and integrable in nature, resulting in the formation of a universal prethermal stage [54–57]. At later times, non-Gaussian correlations develop and integrability is broken, facilitating the transition of the system towards a global equilibrium. While this behavior is well characterized for weakly interacting systems [54], the analogous breaking of integrability for quenches to the strongly interacting regime poses a highly nontrivial theoretical problem due to the appearance of nonperturbative phenomena such as strong few-body scattering, bound states of the medium [57–59], and rapid three-body losses [60]. Furthermore, the quench is associated with the formation of an infinite number of three-body bound Efimov states, whose role and impact remain subjects of active research [59,61–64].

In this Letter, we utilize the method of cumulants to formulate a conserving theory of the quenched unitary Bose gas, capturing the non-Gaussian correlations forming in the *intermediate-time* regime that follows the prethermal stage but preceeds the crossover into a thermal gas. Through direct comparisons with experiment, we characterize the prethermal departure as a consequence of the growth of strong few-body correlations, rather than incoherent particle losses, which broaden the momentum distribution and result in lossless heating of the system. Then, motivated by the absence of a power-law tail in the asymptotics of the momentum distributions observed in Refs. [50,53], we examine how this expected behavior set by the universal contact relations [65–71] is obscured as the system departs the prethermal regime.

Model. We consider a system of N identical bosons of mass m in a cubic volume V, which occupy single-particle states with momentum $\mathbf{k}$ annihilated by operators $\hat{a}_{\mathbf{k}}$. Two such atoms may couple to form an energetically closed channel molecule with center-of-mass momentum $2\mathbf{k}$ and internal energy $\nu > 0$, annihilated by the molecular operator $\hat{b}_{2\mathbf{k}}$. Crucially, we neglect any direct interaction between free atoms, such that all scattering processes in the energetically open channel are mediated by this molecular state. The resulting Hamiltonian reads [72–74]

$$\hat{H} = \sum_{\mathbf{k}} \frac{\hbar^2 k^2}{2m} \hat{a}^\dagger_{\mathbf{k}} \hat{a}_{\mathbf{k}} + \sum_{\mathbf{k}} \left(\frac{\hbar^2 k^2}{4m} + \nu \right) \hat{b}^\dagger_{\mathbf{k}} \hat{b}_{\mathbf{k}} + \frac{g}{2\sqrt{V}} \sum_{\mathbf{k},\mathbf{q}} \left[\zeta(2\mathbf{q})\, \hat{b}^\dagger_{\mathbf{k}} \hat{a}_{\frac{1}{2}\mathbf{k}-\mathbf{q}} \hat{a}_{\frac{1}{2}\mathbf{k}+\mathbf{q}} + \text{H.c.} \right], \quad (1)$$

with $k \equiv |\mathbf{k}|$. The interaction coupling the atomic and molecular states is modeled by a separable potential with strength g and step-function form factor $\zeta(2\mathbf{q}) \equiv \theta(\Lambda - q)$, where θ is the Heaviside step function such that Λ represents a cutoff on the relative two-body momentum. By an analytic solution of the two-body problem, the model parameters can be linked to physical quantities via the renormalization relations [42,75,76],

$$g^2 = \frac{8\pi \hbar^4}{m^2 R_*}, \quad \nu = \frac{\hbar^2}{m R_*} \left(\frac{2\Lambda}{\pi} - a^{-1} \right), \quad (2)$$

where a is the s-wave scattering length, R_* is the characteristic length scale set by the atom-molecule transition rate [77], and the momentum cutoff is related to the van der Waals length as $\Lambda \sim 1/r_{\mathrm{vdW}}$. At unitarity ($a^{-1} \to 0$), the dressed molecular state becomes degenerate with the scattering threshold correspondent with a Feshbach resonance [40]. We express all observables in the Fermi scales $k_n = (6\pi^2 n)^{1/3}$, $E_n = \hbar^2 k_n^2/2m$, and $t_n = \hbar/E_n$, where $n = N/V$ is the gas density. In a many-body context the relative importance of the molecular state is quantified by the resonance width $R_* k_n$ [78].

For both the atomic and molecular fields we adopt the U(1) symmetry-breaking picture of a Bose-Einstein condensate (BEC), where the singlets $\langle a_{\mathbf{k}} \rangle = \delta_{\mathbf{k}0} \sqrt{V} \psi_a$ and $\langle b_{\mathbf{k}} \rangle = \delta_{\mathbf{k}0} \sqrt{V} \psi_m$ are described by the atomic and molecular condensate wave functions ψ_a and ψ_m, respectively. To model the quench scenario, we assume that for $t < 0$ the gas is an ideal atomic BEC, such that $|\psi_a|^2 = n$. At $t = 0$, the system is instantaneously quenched to unitarity such that $na^3 \gg 1$. Then, the far out-of-equilibrium BEC state suffers quantum depletion and sequentially generates higher-order correlations in the gas, which we track using a cumulant expansion [51,79–83]. Truncating the expansion at the level of two-body atomic correlations we obtain the model of Refs. [43,75], with vanishing background scattering length. This *resonance doublet* model includes the cumulants (denoted with subscript c),

$$n^a_{\mathbf{k}} = \langle \hat{a}^\dagger_{\mathbf{k}} \hat{a}_{\mathbf{k}} \rangle_c, \quad \kappa^a_{\mathbf{k}} = \langle \hat{a}_{\mathbf{k}} \hat{a}_{-\mathbf{k}} \rangle_c, \quad (3)$$

which describe the single-particle momentum distribution and pairing field, respectively. Its integrable equations of motion are equivalent to the two-channel Hartree-Fock-Bogoliubov equations [3,43]. In this work, we introduce an extension of this model referred to as the *resonance triplet* model. Here we include the additional molecular and mixed-channel cumulants,

$$n^m_{\mathbf{k}} = \langle \hat{b}^\dagger_{\mathbf{k}} \hat{b}_{\mathbf{k}} \rangle_c, \quad \kappa^m_{\mathbf{k}} = \langle \hat{b}_{\mathbf{k}} \hat{b}_{-\mathbf{k}} \rangle_c,$$
$$\chi_{\mathbf{k}} = \langle \hat{b}^\dagger_{\mathbf{k}} \hat{a}_{\mathbf{k}} \rangle_c, \quad \kappa^{am}_{\mathbf{k}} = \langle \hat{b}_{\mathbf{k}} \hat{a}_{-\mathbf{k}} \rangle_c, \quad (4)$$

and the non-Gaussian tripling field,

$$R^a_{\mathbf{k},\mathbf{q}} = \langle \hat{a}_{\mathbf{k}} \hat{a}_{\mathbf{q}} \hat{a}_{-\mathbf{k}-\mathbf{q}} \rangle_c. \quad (5)$$

In the short-range or vacuum limits, the three-body correlations $\kappa^{am}_{\mathbf{k}}$ and $R^a_{\mathbf{k},\mathbf{q}}$ obey a coupled-channel three-body Schrödinger equation and thus represent closed- and open-channel components of the three-body wave function [76]. As shown in Ref. [59], $R^a_{\mathbf{k},\mathbf{q}}$ becomes macroscopically occupied following the quench and furnishes an order parameter signaling the formation of a Bose-Einstein condensate of Efimovian triples, which motivates its inclusion in the theory. The cumulants $\chi_{\mathbf{k}}$ and $n^m_{\mathbf{k}}$ ensure the conservation of energy, which is possible in the resonance triplet model due to the cubic form of the Hamiltonian [Eq. (1)]. This is a crucial difference between the present model and the single-channel triplet model of Ref. [57], where an inclusion of three-body correlations induces an unphysical leak of energy to the tail of the momentum distribution, thereby prohibiting that model from studying the intermediate-time window considered in this Letter.

Prethermal departure. We begin by examining the dynamics of the momentum distribution in Fig. 1, comparing against experimental results to look for clues about the causes of the prethermal departure and the nature of the intermediate-time regime. Since higher-order correlations develop sequentially following the quench, the doublet and triplet models are initially equivalent [80]. In this early-time regime, atomic pair excitations are directly generated by condensed molecules, as depicted diagrammatically in Fig. 2(a). As shown in Ref. [75], the early-time growth of the condensed molecular fraction scales as $|\psi_m|^2 \sim (t/t_*)^2$, where $t_* = \sqrt{\tau t_n}$ is the mean transition time for atom-molecule conversion. Here $\tau = mR_*/\hbar k_n$ is the molecular lifetime. For broad resonances where $R_* k_n \ll 1$, one finds $t_* \ll t_n$, such that the molecular state acts solely as a mediator of the interaction and the growth of excitations is set purely by t_n. If $R_* k_n$ is increased, the associated increase of t_* gradually slows down the excitation of atoms until the narrow resonance limit $R_* k_n \to \infty$, where all dynamics are frozen.

Following the initial excitation growth, a momentum-dependent plateau is reached, signifying the quasisteady prethermal stage characterized by approximate equilibration of macroscopic observables and emblematic of integrable

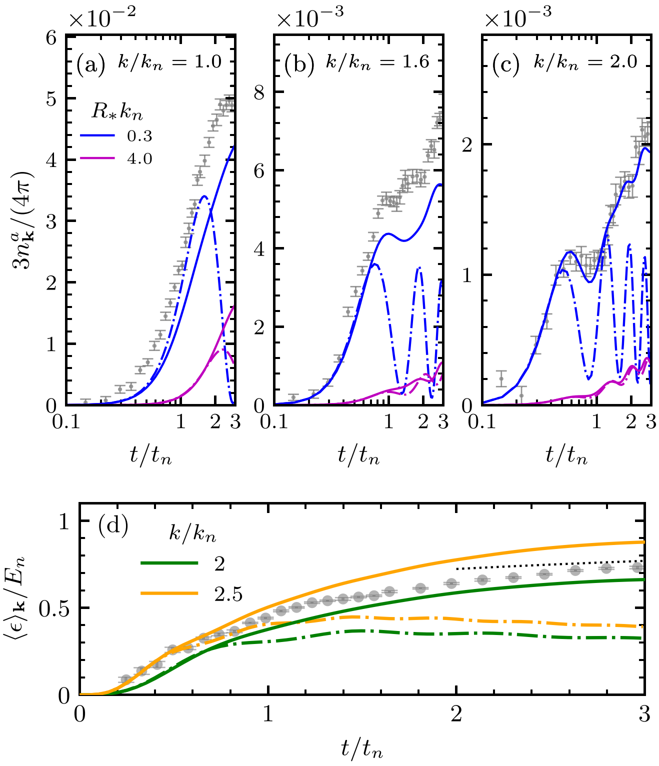

FIG. 1. Broadening of the single-particle momentum distribution following the quench. In panels (a)–(c), we plot the dynamics of the excited-state population $n^a_{\mathbf{k}}$ for a set of momenta. Resonance doublet (triplet) model results are shown with dash-dotted (solid) lines and are matched with the normalization of the experimental data of Ref. [53], shown with gray circles. In panel (d), we show in a similar format the dynamics of the average kinetic energy per particle $\langle\epsilon\rangle_{\mathbf{k}}$ for the broad resonance case $R_* k_n = 0.3$, comparing with the experimental data of Ref. [52]. The black dotted line shows the expected $T \sim t^{2/13}$ scaling for loss-induced heating [52].

dynamics [54,57]. Our key finding is that the rapid depletion of the condensate for broad resonances spurs the growth of R^a and hence of non-Gaussian three-body correlations, leading to a momentum-dependent departure from the prethermal stage visible in Figs. 1(a)–1(c). The shape of the departure shows remarkable qualitative agreement with the experiment of Ref. [53], with just slightly slower growth of excitations due to the mismatch in resonance width. In contrast, the integrable resonance doublet model remains in the prethermal stage at all times, consistent with the single-channel doublet model studied in Ref. [57].

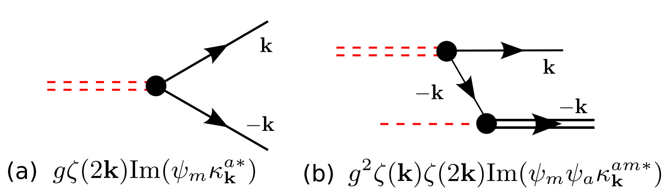

FIG. 2. Two-body (a) and three-body (b) scattering processes that drive the depletion of the atomic condensate. Atomic (molecular) states are shown with single (double) lines, and condensed states are shown with dashed red lines. Each vertex represents atom-molecule conversion.

Physically, the broadening of the momentum distribution shown in Fig. 1 arises from the significant interaction energy injected into the system by the quench, which is gradually converted into kinetic energy via two- and three-body scattering and thus results in lossless correlation-induced, rather than recombination-induced, heating. To quantify this lossless heating, we follow experiment [52] and examine the averaged kinetic energy per particle [57],

$$\langle\epsilon\rangle_{\mathbf{k}} = \frac{1}{N}\sum_{\mathbf{k}'} n^a_{\mathbf{k}'}\frac{\hbar^2 k'^2}{2m}\theta(k-k'). \tag{6}$$

The restriction to a maximum momentum k accounts for the limited resolution in the experiment [52]. From Fig. 1(d) it is clear that the Gaussian statistics of the resonance doublet model are unable to capture the experimental departure from the prethermal plateau that occurs by $t \sim t_n$. The resonance triplet model does capture this departure, following the experimental results until the crossover where the gas was experimentally found to pass from degenerate to nondegenerate regimes. This agreement further strengthens the interpretation of the dynamics at intermediate times as correlation (rather than loss) dominated.

Correlations in the intermediate-time regime. Having provided evidence for the lossless nature of the intermediate-time regime, we now study the dominant processes in the system during this time. The difference between the resonance doublet and resonance triplet dynamics in Fig. 1 arises predominantly from a distinct non-Gaussian process in the resonance triplet model, shown diagrammatically in Fig. 2(b) [76]. Together, the processes in Fig. 2 seed the postquench growth of short-range atomic two- and three-body correlation functions, given as

$$\langle d^\dagger d\rangle = \left|\frac{1}{V}\sum_{\mathbf{k}}\kappa^a_{\mathbf{k}}\right|^2, \quad \langle t^\dagger t\rangle = \left|\frac{1}{V^{\frac{3}{2}}}\sum_{\mathbf{k},\mathbf{q}} R^a_{\mathbf{k},\mathbf{q}}\right|^2. \tag{7}$$

Here $\hat{d} = \hat{\psi}(\mathbf{0})\hat{\psi}(\mathbf{0})$ and $\hat{t} = \hat{\psi}(\mathbf{0})\hat{\psi}(\mathbf{0})\hat{\psi}(\mathbf{0})$ represent the local dimer and trimer field, respectively, with the field operator $\hat{\psi}(\mathbf{0}) = (1/\sqrt{V})\sum_{\mathbf{k}}\hat{a}_{\mathbf{k}}$. Due to their short-range nature, the value of the correlation functions is expected to be fully specified by universal relations derived from few-body physics. Hence they can be directly related to the two-body and three-body contact densities $\mathcal{C}_2$ and $\mathcal{C}_3$, which form a measure of the probability for finding two and three particles in close proximity [65–67,70]. In our two-channel model, we derive the following relations from effective field theory [76],

$$\mathcal{C}_2 = \frac{m^2 g^4}{4\hbar^4\nu^2}\langle d^\dagger d\rangle + \frac{m^3 g^6}{2\hbar^6\nu^3\Lambda^2}\left(H + \frac{J}{\pi}\right)\langle t^\dagger t\rangle, \tag{8}$$

$$\mathcal{C}_3 = -\frac{m^2 g^4}{8\hbar^4\nu^2\Lambda^2}H'\langle t^\dagger t\rangle. \tag{9}$$

Here H, J, and H' are known log-periodic functions of the Efimovian binding wave number κ_* and the cutoff Λ [70]. As shown in Fig. 3, the prethermal departure is indeed correlated with a significant increase in $\mathcal{C}_3$, indicating the introduction of strong non-Gaussian three-body correlations in this

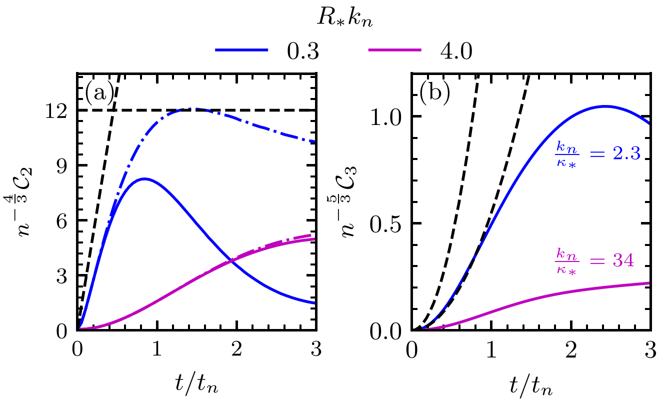

FIG. 3. Dynamics of two-body and three-body contact densities $\mathcal{C}_2$ in panel (a) and $\mathcal{C}_3$ in panel (b), for different values of the resonance width $R_* k_n$. Results from the resonance doublet (triplet) model are shown with dash-dotted (solid) lines. For comparison we show with black dashed lines the linear early-time growth of $\mathcal{C}_2$ as derived in Ref. [84] for broad resonances and the long-time constant value of $\mathcal{C}_2$ from Ref. [85]. For $\mathcal{C}_3$ we show the range of quadratic early-time growths derived in Ref. [63], which correspond to maximum ($k_n/\kappa_* \sim 1$) and minimum ($\kappa_* \ll k_n$) enhancement of $\mathcal{C}_3$ due to the Efimov effect, again in the broad resonance limit.

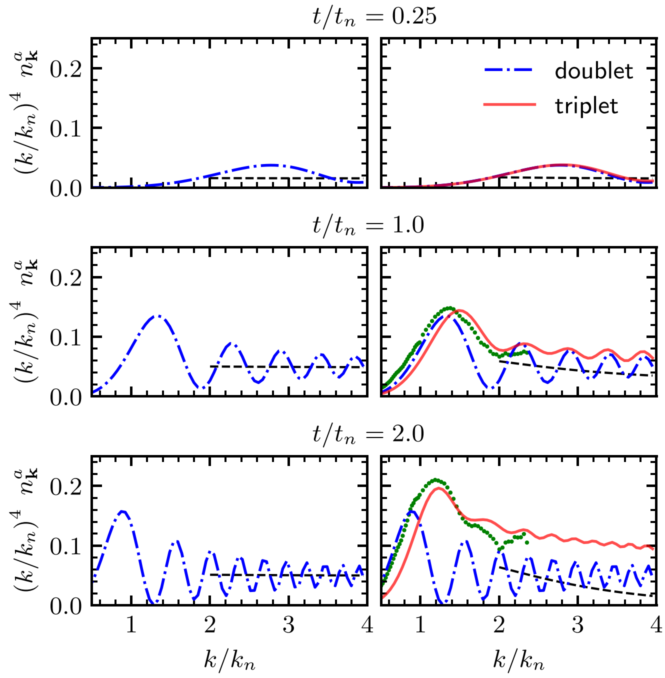

FIG. 4. Dynamics of the tail of the single-particle momentum distribution for the broad resonance $R_* k_n = 0.3$. In the left-hand (right-hand) panels, we compare the resonance doublet (triplet) models with the asymptotic prediction in Eq. (10), shown with a black dashed line. For the sake of comparison the doublet results are replotted in the right-hand panels. For $t/t_n = 1.0$ and 2.0, we also compare with the experimental data of Ref. [53], shown with the green circles.

regime. Simultaneously, the influence of surrounding particles on clustered pairs leads to a significant decrease of the two-body contact [57]. As shown in Ref. [59], the value of $\mathcal{C}_2$ and $\mathcal{C}_3$ following the quench is determined predominantly by the magnitude of macroscopic order parameters associated with condensation of pairs and triples. These findings are confirmed in our model by the development of a substantial fraction of condensed triples for $t > t_n$ [76]. We note that the increase of $\mathcal{C}_3$ for broad resonances is connected with an increased overlap of the size of the Efimov trimer, quantified by the binding wave number κ_*, and the Fermi scale k_n. Consistent with Figs. 1(a)–1(c), the value of $\mathcal{C}_3$ is decreased for narrow resonances, where $\kappa_* \ll k_n$ represents the previously unexplored limit opposite to the universal regime $\kappa_* \gg k_n$ examined in Ref. [59].

Asymptotics in the intermediate-time regime. We now focus on the large-momentum tail of $n^a_{\mathbf{k}}$, whose expected behavior due to universal relations derived from few-body physics was not observed in experiments [50,53]. Specifically, the following asymptotic behavior is predicted at thermal equilibrium,

$$k^4 n^a_{\mathbf{k}} \underset{k\to\infty}{\longrightarrow} \mathcal{C}_2 + \mathcal{C}_3 F(k)/k, \tag{10}$$

where $F(k)$ is a log-periodic function specified by κ_* [70]. While Eq. (10) provides a possible route to extract the values of the contacts from $n^a_{\mathbf{k}}$, a $1/k^4$ tail has so far not been observed in quench experiments [50,53]. Fits to Eq. (10) were made in Ref. [71], but found values for $\mathcal{C}_2$ were considerably larger than expected from theoretical calculations [57,85,86]. Additionally it has been shown that in lower-dimensional systems the scaling in Eq. (10) can be disrupted by nonlocal correlations [87].

Motivated by this disagreement between existing theory and experiment, we compare in Fig. 4 the single-particle momentum distribution obtained from our models with the asymptotic prediction in Eq. (10). In the resonance doublet model, where all dynamics are Gaussian, $\mathcal{C}_3$ vanishes trivially and one can show analytically that the expected $1/k^4$ power law due to $\mathcal{C}_2$ is always present [76]. In the resonance triplet model, we observe that the non-Gaussian processes responsible for the departure from the prethermal stage for $t \gtrsim t_n$ damp the oscillatory behavior of $n^a_{\mathbf{k}}$, consistent with the findings of Ref. [54] at weak interactions and in agreement with the experimental data of Ref. [53]. At the same time, the significant increase of excitations for $k > k_n$ obscures the expected power laws for all the examined momenta $k/k_n \leqslant 4$. Hence, once the gas has exited the prethermal stage, the asymptotic expansion in Eq. (10) no longer captures the momentum distribution over the considered range. As experimental results are lacking beyond this range due to poor signal-to-noise ratios [50], our findings suggest that significant caution should be exercised when fitting Eq. (10) out of equilibrium and explain why a power-law tail was not seen in Refs. [50,71]. It is important to note, however, that our results do not invalidate Eq. (10), but rather push its possible applicability to larger momenta. For sufficiently large values of k/k_n, our distributions do converge consistently with the value of $\mathcal{C}_2$ obtained from Eq. (8), but also exhibit strong finite-range features due to the density regimes considered [76]. We have confirmed that the distributions in Fig. 4 are density independent.

Conclusion. In this Letter, we have shown how a conserving many-body model constructed from a selection of microscopically relevant Gaussian and non-Gaussian correlations is able to elucidate the dynamics of quenched unitary Bose gases in the time window succeeding the integrable dynamics associated with the prethermal stage. In the future, this general framework can be used in studying the wide array of other quantum systems that exhibit non-Gaussian physics, including, for example, strongly interacting ultracold mixtures and polarons [29,48,49,88,89], trions in semiconductors [90–94], nuclear matter [95–98], and Rydberg atom arrays [99].

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