Third- and Fourth-Order Virial Coefficients of Harmonically Trapped Fermions in a Semiclassical Approximation

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Third- and fourth-order virial coefficients of harmonically trapped fermions in a semiclassical approximation

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Using a leading-order semiclassical approximation, we calculate the third- and fourth-order virial coefficients of nonrelativistic spin-1/2 fermions in a harmonic trapping potential in arbitrary spatial dimensions, and as functions of temperature, trapping frequency, and coupling strength. Our simple, analytic results for the interaction-induced changes \( \Delta b_3 \) and \( \Delta b_4 \) agree qualitatively, and in some regimes quantitatively, with previous numerical calculations for the unitary limit of three-dimensional Fermi gases.

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I. INTRODUCTION

The properties of fermions at finite temperature and density are in part governed by the dimensionless product \( \beta \mu \), where \( \beta = 1/(k_B T) \) is the inverse temperature and \( \mu \) is the chemical potential. Typically, the region \( \beta \mu \approx 0 \) displays a crossover between quantum and classical physics, while \( z = e^{\beta \mu} \ll 1 \) indicates a dilute limit where the thermodynamics is given by the virial expansion, which expands a given physical quantity as functions of temperature, trapping frequency, and coupling strength. Our simple, analytic results for the interaction-induced changes \( \Delta b_3 \) and \( \Delta b_4 \) agree qualitatively, and in some regimes quantitatively, with previous numerical calculations for the unitary limit of three-dimensional Fermi gases.

II. HAMILTONIAN AND FORMALISM

As our focus is on systems with short-range interactions, such as dilute atomic gases or dilute neutron matter, the Hamiltonian reads

\[
\hat{H} = \hat{H}_0 + \hat{V}_{\text{int}},
\]

where

\[
\hat{H}_0 = \hat{T} + \hat{V}_{\text{ext}},
\]

and

\[
\hat{T} = \sum_{j=1,2} \int d^d x \, \hat{\psi}_j^\dagger(x) \left( -\frac{\hbar^2}{2m} \nabla^2 \right) \hat{\psi}_j(x)
\]

is the kinetic energy.

\[
\hat{V}_{\text{ext}} = \frac{1}{2} m \omega^2 \int d^d x \, x^2 [\hat{\rho}_1(x) + \hat{\rho}_2(x)],
\]

is the spherically symmetric external trapping potential, and

\[
\hat{V}_{\text{int}} = -g_d \int d^d x \, \hat{\rho}_1(x) \hat{\rho}_2(x),
\]

is the interaction.

In the above equations, the field operators \( \hat{\psi}_j, \hat{\psi}_j^\dagger \) correspond to particles of species \( s = 1, 2 \), and \( \hat{\rho}_s(x) \) are the coordinate-space densities. For the remainder of this work, we will set \( h = k_B = m = 1 \).

A. Thermodynamics and the virial expansion

The equilibrium thermodynamics of our quantum many-body system can be captured by the grand-canonical partition function, namely,

\[
\mathcal{Z} = \text{tr}[e^{-\beta(\hat{H} - \mu \hat{N})}] = e^{-\beta \Omega},
\]

where \( \beta \) is the inverse temperature, \( \Omega \) is the grand thermodynamic potential, \( \hat{N} \) is the total particle number operator, and \( \mu \) is the overall chemical potential (we will not consider polarized systems in this work).

As the calculation of \( \mathcal{Z} \) is a formidable problem in the presence of interactions, we resort to approximations or numerical evaluations in order to access the thermodynamics. To that end, in this work we will use the virial expansion, which is an expansion around the dilute limit \( z \rightarrow 0 \), where \( z = e^{\beta \mu} \) is the fugacity, i.e., it is a low-fugacity expansion (see Ref. [4] for a review on recent applications of the virial expansion to ultracold atoms). The coefficients accompanying the powers of \( z \) in the expansion \( \Omega \) are the virial coefficients \( b_n \),

\[
-\beta \Omega = \ln \mathcal{Z} = Q_1 \sum_{n=1}^\infty b_n z^n,
\]
where \( Q_1 \) is the one-body partition function. Using the fact that \( Z \) is itself a sum over canonical partition functions \( Q_N \) of all possible particle numbers \( N \), namely,

\[
Z = \sum_{N=0}^{\infty} z^N Q_N,
\]

we obtain expressions for the virial coefficients

\[
b_1 = 1, \quad b_2 = \frac{Q_2}{Q_1} - \frac{Q_1}{2!}, \quad b_3 = \frac{Q_3}{Q_1} - b_2 Q_1 - \frac{Q_1^2}{3!}, \quad b_4 = \frac{Q_4}{Q_1} - \left( b_3 + \frac{b_2^2}{2} \right) Q_1 - b_2 Q_1^2 - \frac{Q_1^3}{4!},
\]

and so on. In this work we will not pursue the virial expansion beyond \( b_4 \). The \( Q_N \) can themselves be written in terms of the partition functions \( Q_{a,b} \) for \( a \) particles of type 1 and \( b \) particles of type 2,

\[
Q_1 = 2Q_{1,0}, \quad Q_2 = 2Q_{2,0} + Q_{1,1}, \quad Q_3 = 2Q_{3,0} + 2Q_{2,1}, \quad Q_4 = 2Q_{4,0} + 2Q_{3,1} + Q_{2,2},
\]

and so on for higher orders. In the absence of intraspecies interactions, only \( Q_{1,1}, Q_{2,1}, Q_{3,1}, \) and \( Q_{2,2} \) are affected, such that the change in \( b_2, b_3, \) and \( b_4 \) due to interactions is entirely given by

\[
\Delta b_2 = \frac{\Delta Q_{1,1}}{Q_1}, \quad \Delta b_3 = \frac{2\Delta Q_{2,1}}{Q_1} - \Delta b_2 Q_1, \quad \Delta b_4 = \frac{2\Delta Q_{3,1} + \Delta Q_{2,2}}{Q_1} - \Delta \left( b_3 + \frac{b_2^2}{2} \right) Q_1 - \Delta b_2 Q_1^2 - \frac{Q_1^3}{4!}.
\]

To calculate \( \Delta Q_{m,n} \), we implement a semiclassical approximation, as described in the next section. Once we obtain the virial coefficients, one may rebuild the grand-canonical potential \( \Omega \) to access the thermodynamics of the system as a function of the various parameters.

In order to connect to the physical parameters of the systems at hand, we will use the value of \( \Delta b_2 \) as a renormalization condition by relying on the exact answers as functions of \( \beta \omega \) and the physical coupling \( \lambda \). These exact answers are not always known analytically, but they can easily be obtained numerically by solving the two-body problem of interest.

Although in this work we will focus on systems in a harmonic trap, thus far the identities presented in this section are more general. As a reference for the trapped case, we present here the calculation of the noninteracting virial coefficients for arbitrary \( \beta \omega \). (We note that such a calculation, while simple, does not appear in the literature.) Starting from the logarithm of the noninteracting partition function in \( d \) spatial dimensions, we have, for two fermion species,

\[
\ln Z = 2 \sum_n \ln \left( 1 + z e^{-\beta \omega d/2} \prod_{i=1}^{d} e^{-\beta \omega n_i} \right).
\]

Expanding in powers of \( z \) on both sides, and switching the order of the sums, we obtain

\[
Q_1 \sum_{k=1}^{\infty} b_1^k z^k = 2 \sum_{k=1}^{\infty} \frac{(-1)^{k+1}}{k} e^{-\beta \omega d k/2} \left( \sum_{n=0}^{\infty} e^{-\beta \omega n} \right)^d = 2 \sum_{k=1}^{\infty} \frac{(-1)^{k+1}}{2^k k} \left( \frac{1}{\sinh(\beta \omega k/2)} \right)^d.
\]

To identify the noninteracting virial coefficients \( b_n^0 \), we need \( Q_1 \),

\[
Q_1 = 2 \sum_n e^{-\beta E_n} = 2 e^{-\beta \omega d/2} \left( \frac{1}{1 - e^{-\beta \omega}} \right)^d \quad \text{(22)}
\]

Thus, the virial coefficients of a trapped noninteracting spin-1/2 Fermi gas in \( d \) dimensions are

\[
b_n^0 = \frac{(-1)^{n+1}}{n} \left( \frac{\sinh(\beta \omega/2)}{\sinh(\beta \omega n/2)} \right)^d. \quad \text{(24)}
\]

Notably, in the limit \( \beta \omega \ll 1 \), we obtain

\[
b_n^0 \rightarrow (-1)^{n+1} \left( \frac{1}{n} \right)^{d+1},
\]

which agrees in \( d = 3 \) with the local density approximation result quoted in Ref. [4]. The simple result of Eq. (24) should be a textbook calculation, but it does not appear elsewhere, to the best of our knowledge. Note that for the homogeneous (i.e., untrapped) system, the noninteracting virial coefficients in \( d \) dimensions are

\[
b_n^0 \rightarrow \frac{(-1)^{n+1}}{n} \left( \frac{1}{n} \right)^{d+1},
\]

such that \( b_n^0 = b_n^0 \hom n^{-d} \) for \( \beta \omega \ll 1 \).

B. semiclassical lattice approximation

To calculate the interaction-induced change \( \Delta Q_{m,n} \), we implement an approximation which consists in keeping the leading term in the commutator expansion,

\[
e^{-\beta (\hat{H}_0 + \hat{V}_{\text{int}})} = e^{-\beta \hat{H}_0} e^{-\beta \hat{V}_{\text{int}}} \times e^{-\frac{\hat{V}_{\text{int}}}{(\hat{H}_0, \hat{V}_{\text{int}})}} \times \ldots,
\]

where the higher orders involve exponentials of nested commutators of \( \hat{H}_0 \) with \( \hat{V}_{\text{int}} \). Thus, the leading order in this expansion consists in setting \( [\hat{H}_0, \hat{V}_{\text{int}}] = 0 \), which corresponds to a semiclassical approximation. Another way to see this approximation is in terms of a Trotter-Suzuki factorization, i.e.,

\[
e^{-\beta (\hat{H}_0 + \hat{V}_{\text{int}})} \rightarrow \lim_{n \to \infty} e^{-\beta \hat{H}_0/2} e^{-\beta \hat{V}_{\text{int}}/2n} \ldots,
\]

where the \( n = 1 \) case is the leading order we pursue in this work and higher orders can be defined by increasing \( n \).
C. Example: Calculation of $\Delta Q_{1,1}$ and $\Delta b_2$

In the approximation proposed above, the two-particle problem is analyzed as follows:

$$Q_{1,1} = \text{Tr}[e^{-\beta H_0} e^{-\beta \hat{V}_{\text{int}}}]$$

$$= \sum_{k_1, k_2, x_1, x_2} \langle k_1 | x_1, x_2 \rangle e^{-\beta(k_1 + k_2)} \langle x_1, x_2 | e^{-\beta \hat{V}_{\text{int}}} | k_2, x_2 \rangle$$

$$= \sum_{k_1, k_2, x_1, x_2} e^{-\beta(E_{k_1} + E_{k_2})} M_{x_1, x_2} | \langle k_1 | k_2 | x_1 | x_2 \rangle |^2, \quad (29)$$

where we have inserted complete sets of states in coordinate space $\{|x_1, x_2\rangle\}$ and in the basis $|k_1, k_2\rangle$ of eigenstates of $\hat{H}_0$, whose single-particle eigenstates $|k\rangle$ have eigenvalues $E_k$. We have also made use of the fact that $\hat{V}_{\text{int}}$ is diagonal in coordinate space, such that

$$M_{x_1, x_2} = 1 + C \delta_{x_1, x_2}, \quad (30)$$

where

$$C = (e^{\beta \delta} - 1) \ell^d. \quad (31)$$

Here [5], we treated the interaction by placing the system on a spatial lattice of spacing $\ell$, as in Ref. [2], which will set the scale for all the dimensionful quantities from this point on. We therefore assume from this point on that any remaining powers of $\ell$ have been absorbed into $g_d$, which now represents a lattice coupling, such that the combination $\beta g_d$ is dimensionless.

Thus,

$$\Delta Q_{1,1} = C \sum_{k_1, k_2, x} \ell^d e^{-\beta(E_{k_1} + E_{k_2})} |\phi_k(x)\rangle^2 |\phi_k(x)\rangle^2, \quad (32)$$

and we will use normalized single-particle wave functions in Cartesian coordinates which in one dimension (1D) take the form

$$\phi_k(x) = \frac{1}{\sqrt{2\pi \hbar}} \left( \frac{\omega}{\pi} \right)^{1/4} e^{-\omega x^2/2} H_n(\sqrt{\omega} x), \quad (33)$$

where the $H_n$ are Hermite polynomials. We note in passing that, in $d$ dimensions, $\phi_k(x)$ has units of $\ell^{-d/2}$, since $\omega$ has units of $\ell^{-2}$, such that $\Delta Q_{1,1}$ is dimensionless as expected.

The sums over $k_1$ and $k_2$ in Eq. (32) are independent and identical and take the form

$$\sum_{n=0}^{\infty} e^{-\beta E_n} |\phi_n(x)\rangle^2 = \frac{\sqrt{\omega}}{\pi} e^{-\beta \omega/2 - \omega x^2} G(\beta \omega, \sqrt{\omega} x), \quad (34)$$

for each Cartesian dimension, where the function $G$ can be calculated as a special case of Mehler’s formula [6],

$$G(k, y) \equiv \sum_{n=0}^{\infty} \frac{e^{-kn}}{2\pi \hbar} [H_n(y)]^2 = \frac{\exp[2y^2/(1 + e^{2y})]}{\sqrt{1 - e^{-2y}}}. \quad (35)$$

This formula encodes the finite-temperature, single-particle density matrix of a noninteracting, nonrelativistic system in a harmonic trapping potential, and therefore its use is essential in the calculations that follow.

Squaring the result, we obtain, in one spatial dimension,

$$\Delta Q_{11} = 2C \int_0^\infty dx \frac{\omega}{2\pi \sinh(\beta \omega)} \exp[-2\omega x^2 \tanh(\beta \omega/2)]$$

$$= C \frac{\sqrt{\omega}}{4} \frac{1}{\sqrt{\pi \sinh(\beta \omega) \sinh(\beta \omega/2)}}, \quad (35)$$

where we have performed the last Gaussian integral along with some hyperbolic function simplifications. Note that $\Delta Q_{11}$ is a dimensionless quantity; we have, in the above equation, implicitly replaced a sum with an integral by taking the continuum limit. In doing so, notice that $C$ has dimensions of $\ell^d$ [see Eqs. (30) and (31), and formulas below].

Generalizing to $d$ spatial dimensions,

$$\Delta Q_{11} = C \left[ \frac{\sqrt{\omega}}{4} \frac{1}{\sqrt{\pi \sinh(\beta \omega) \sinh(\beta \omega/2)}} \right]^d \quad (36)$$

Using Eq. (22), we find that $Q_1$ cancels exactly in the final expression, as expected, such that

$$\Delta b_2 = C \left[ \frac{\sqrt{\omega}}{2\sqrt{\pi \sinh(\beta \omega)}} \right]^d = C \frac{1}{2 \sinh(\beta \omega)} \left[ \frac{2\sinh(\beta \omega)}{\beta \omega} \right]^2 \quad (37)$$

where we have used the thermal wavelength $\lambda_T = \sqrt{2\pi \beta}$ to write the result in as a function of $\beta \omega$.

As mentioned above, we will use this result to connect to the physical coupling $\lambda - \lambda_1$ of a given system, as a renormalization condition at a given value of $\beta \omega$. To that end, we first solve for $C/\lambda_T^d$,

$$\frac{C}{\lambda_T^d} = 2\Delta b_2 \left[ \frac{2\sinh(\beta \omega)}{\beta \omega} \right]^2 \quad (38)$$

In the unitary limit of the 3D Fermi gas [7], for instance, the exact answer for $\Delta b_2$ is known [4],

$$\Delta b_2 = \frac{1}{2} \left( \frac{e^{-\beta \omega/2}}{1 + e^{-\beta \omega}} \right) = \frac{1}{4 \cosh(\beta \omega/2)} \quad (39)$$

Using that result in Eq. (38) yields

$$\frac{C}{\lambda_T^d} = e^{-\beta \omega/2} \left[ \frac{2\sinh(\beta \omega)}{\beta \omega} \right]^{3/2} \quad (40)$$

As we will show below, this type of renormalization is very practical as our results for $\Delta b_2$ are simple quadratic functions of $C/\lambda_T^d$ with $\beta \omega$-dependent coefficients.

III. RESULTS

Following the steps outlined above in the example calculation of $\Delta b_2$, we have calculated the various contributions to $\Delta b_3$ and $\Delta b_4$, which we present in this section. In all cases, the central component of the calculation is the use of the analytic form of Mehler’s kernel, which effectively reduces the calculation to a small number of Gaussian integrals.

A. Result for $\Delta Q_{2,1}$ and $\Delta b_3$

With small modifications to the example for $\Delta b_2$, it is straightforward to show that

$$\Delta Q_{2,1} = C \lambda_T^d \left[ \frac{\beta \omega}{4 \sinh^3(\beta \omega)} \right]^2 \times \left[ \frac{1}{2^2 \tanh^2(\beta \omega/2)} - \left( \frac{2 \cosh^2(\beta \omega/2)}{4 \cosh^2(\beta \omega/2) - 1} \right)^4 \right] \quad (41)$$
and using the results of the previous section, it is easy to assemble the final answer for $\Delta b_3$ in our approximation using

$$\Delta b_3 = \frac{2\Delta Q_{2,1}}{Q_1} - \Delta b_2 Q_1.$$  \hfill (42)

Note that $Q_1$ diverges in the $\beta \omega \to 0$ limit, but that divergence will cancel out in the final expression for $\Delta b_3$. Indeed, after simplifications, we obtain

$$\Delta b_3 = -\frac{C}{\lambda_f^2} \left[ \frac{\beta \omega}{2 \sinh(\beta \omega)} \right]^d \left( \frac{1}{4 \cosh^2(\beta \omega/2) - 1} \right)^{\frac{d}{2}} \frac{1}{2 \sinh^2(\beta \omega)} \right]^{2d} \left( \frac{1}{2 \sinh^2(\beta \omega)} - 1 \right)^{\frac{d}{2}}. \hfill (43)$$

The final answer for $\Delta b_3$ in terms of $\Delta Q_{2,1}$ and $\Delta b_2$ can be assembled using

$$\Delta b_3 = \frac{2\Delta Q_{2,1}}{Q_1} + \Delta \left( b_1 + \frac{b_2}{2} \right) Q_1 - \Delta b_2 Q_1.$$  \hfill (48)

After several simplifications and cancellations (which can be tracked by their degree of divergence as $\beta \omega \to 0$, while the final result for $\Delta b_3$ is finite), we obtain

$$\Delta b_3 = \frac{C}{\lambda_f^2} \left[ \frac{\sqrt{2\beta \omega \sinh(\beta \omega/2)}}{2 \sinh^2(\beta \omega)} \right]^d \left\{ \frac{1}{2 \tanh^2(\beta \omega/2)} + \frac{2 \coth^2(\beta \omega/2)}{2 \coth(\beta \omega/2)} \right\}. \hfill (49)$$

In this case, the limit $\beta \omega \to 0$ yields

$$\Delta b_4 \to \frac{C}{\lambda_f^2} 4^{-\frac{d}{2}} (3^{-\frac{d}{2}} + 2^{-d-1}) + \left( \frac{C}{\lambda_f^2} \right)^2 4^{-d-1} (1 - 2^{d+1}).$$  \hfill (50)

Once again, we recall the homogeneous result

$$\Delta b_4^{\text{hom}} = \frac{C}{\lambda_f^2} 3^{-\frac{d}{2}} + 2^{-d-1} + \left( \frac{C}{\lambda_f^2} \right)^2 4^{-d-1} (1 - 2^{d+1}).$$  \hfill (51)

which is manifestly finite in the $\beta \omega \to 0$ limit. In that limit,

$$\Delta b_3 \to -\frac{C}{\lambda_f^2} \frac{1}{6\beta \omega}. \hfill (44)$$

We recall the result of Ref. [1] for the homogeneous case, namely,

$$\Delta b_3^{\text{hom}} = -\frac{C}{\lambda_f^2} \frac{1}{2\beta \omega}, \hfill (45)$$

which shows that the relationship between the homogeneous and trapped cases, pointed out in the Introduction, is also satisfied once interactions are turned on.

### B. Result for $\Delta Q_{3,1}, \Delta Q_{2,2},$ and $\Delta b_4$

Again following the steps outlined above, we obtain

$$\Delta Q_{2,2} = \frac{C}{\lambda_f^2} \left[ \frac{\sqrt{2\beta \omega \sinh(\beta \omega/2)}}{2 \sinh^2(\beta \omega)} \right]^d \left\{ \frac{1}{2 \tanh^2(\beta \omega/2)} + \frac{2 \coth^2(\beta \omega/2)}{2 \coth(\beta \omega/2)} \right\}. \hfill (46)$$

and find that $\Delta b_4 = \Delta b_4^{\text{hom}} 4^{-\frac{d}{2}}$, as expected.

### C. Results in terms of $\Delta b_2$

Finally, collecting our results for $\Delta b_1$ and $\Delta b_2$ and expressing them in terms of $\Delta b_2$ [via Eq. (38)], we obtain

$$\Delta b_3 = -2 \left( \frac{1}{4 \cosh^2(\beta \omega/2) - 1} \right)^{\frac{d}{2}} \Delta b_2.$$  \hfill (52)

and

$$\Delta b_4 = \left( \frac{1}{2 \cosh(\beta \omega/2)} \right)^{\frac{d}{2}} f_1(\beta \omega, d) \Delta b_2 + f_2(\beta \omega, d) (\Delta b_2)^2,$$  \hfill (53)

where

$$f_1(\beta \omega, d) = \left( \frac{1}{2 \cosh(\beta \omega/2)} \right)^{\frac{d}{2}} + 2 \left( \frac{1}{4 \cosh^2(\beta \omega/2) - 1} \right)^{\frac{d}{2}},$$  \hfill (54)

and

$$f_2(\beta \omega, d) = \left( \frac{1}{2 \cosh(\beta \omega/2)} \right)^{\frac{d}{2}} - 2 \left( \frac{1}{4 \cosh^2(\beta \omega/2) - 2} \right)^{\frac{d}{2}}.$$  \hfill (55)

The above formulas for $\Delta b_3$ and $\Delta b_4$ are the main result of this work. In the following, we explore their behavior as a function of $d$ and $\beta \omega$, focusing in particular on the unitary limit of the 3D Fermi gas. While numerical results exist for
these quantities in some cases, in particular in 2D [8–11] (see also Refs. [12–14]) and in 3D at unitarity [15–19], most of those correspond to homogeneous systems and do not feature explicit, analytic dependence on the dimension nor on $\beta\omega$, as shown here. Our results are therefore useful in that they are able to provide analytic insight into the behavior of virial coefficients across dimensions, and as a function of the temperature (or trapping frequency) as well as the coupling strength. Below, we evaluate our formulas and discuss the resulting answers.

**D. Qualitative behavior**

To illustrate our analytic results, in Fig. 1 we show $\Delta b_3/\Delta b_2$ (top) and $\Delta b_4/\Delta b_2$ (bottom) as a function of the spatial dimension $d$, at various $\beta\omega$, fixing $\Delta b_2$ to its value in the unitary limit (as a reference point). We find that, as $d$ increases, the magnitude of the interaction-induced change $\Delta b_n$ decreases. This suggests that, using $\Delta b_2$ as the fixed, dimension-independent coupling, the radius of convergence of the virial expansion increases with $d$. This is consistent with the idea that, in higher dimensions, the kinetic energy dominates over the interactions and mean-field type of approaches capture the behavior of the system correctly.

As a comparison with previous calculations, we show in Fig. 2 our results in 3D at unitarity as a function of $\beta\omega$, superimposed with the data from Ref. [20]. While we do not expect, a priori, good quantitative agreement in this strong-coupling regime, we find at least qualitative agreement for both $\Delta b_3$ and $\Delta b_4$, and surprisingly good agreement at the quantitative level for $\Delta b_4$. Clearly, the LO-SCLA is able to capture more than just the shape of the $\beta\omega$ dependence of the virial coefficients.

**IV. SUMMARY AND CONCLUSIONS**

In this article we have implemented a semiclassical approximation, at leading order, to calculate the virial coefficients $\Delta b_3$ and $\Delta b_4$ of harmonically trapped Fermi gases. Our calculations yield analytic answers as functions of $\beta\omega$ and, by defining the lattice coupling $C$ so as to match $\Delta b_2$ to the known exact result, we also obtain the dependence on the physical coupling strength. Notably, our results are also analytic functions of the spatial dimension $d$, allowing us to study the behavior of the virial expansion across dimensions. We find that, at fixed $\Delta b_2$, the magnitude of $\Delta b_n$ decreases as $d$ increases, for all $\beta\omega$.

Although there have been many (and very precise) determinations of virial coefficients in the literature, they are mostly numerical and focus on specific dimensions or couplings (and most of them are for homogeneous systems). Our approach...
and results are, in that sense, complementary: We do not expect high precision from the LO-SCLA, but through it we are able to study, explicitly, the variations with the parameters of the problem, which yield qualitative analytic insight into the properties of the virial expansion. We have demonstrated the quality of our leading-order results for $\Delta b_3$ and $\Delta b_4$ in the unitary limit by showing that they qualitatively follow the expected answers, which is an encouraging sign to proceed to next-to-leading order in future work. Furthermore, the approximate agreement with prior results at unitarity suggests that, between the noninteracting regime and the unitary point, that agreement should be even better than shown here.

We would like to stress that the calculation presented here is nonperturbative in the same sense as in lattice Monte Carlo approaches, except that we have implemented a very coarse temporal lattice. Increasing the order of the approximation (beyond LO-SCLA studied here) amounts to reducing the temporal lattice spacing effects. Those effects disappear in the limit where either $\hat{H}_0$ or $\hat{V}_{\text{int}}$ appear on their own. When both are present, the regime of validity of the SCLA at a given order is controlled by the magnitude of the commutator $[\hat{H}_0, \hat{V}_{\text{int}}]$, i.e., it will depend on the coupling strength.

Finally, it should be pointed out that the renormalization of the lattice theory based on $\Delta b_2$ does not by itself eliminate all the lattice artifacts. Future studies should explore the use of improved actions (see, e.g., Refs. [22,23]), potentially making use of prior knowledge of $\Delta b_3$ where available, to enhance the quality of the expansion.

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