

Temperature Dependence of p -Wave Contacts in a Harmonically Trapped Fermi Gas

Kenta Nagase¹, Hikaru Takahashi¹, Soki Oshima¹, and Takashi Mukaiyama^{1*}

Department of Physics, Institute of Science Tokyo, Ookayama 2-12-1, Meguro-ku, Tokyo 152-8550, Japan



(Received 4 September 2025; accepted 2 December 2025; published 8 January 2026)

We study the dependence of the p -wave contact on the Fermi temperature T_F and reduced temperature T/T_F based on the number of closed-channel molecules. From the anisotropic pattern of dissociated molecules, we resolve the narrow $m = 0$ and $m = \pm 1$ dipolar splitting of the p -wave Feshbach resonance in ^6Li , enabling the independent determination of the contact for all three m components. For each component, we identify a near-resonant scaling with $\sqrt{T_F}$, indicating the contribution of the normalized effective range $k_F R_e$. In addition, we show that the peak contacts observed near resonance increase as T/T_F is lowered, a trend that is accurately captured by estimates based on the second-order virial expansion. Our results, together with estimates from the p -wave virial expansion, provide a route toward a complete understanding of the thermodynamics of resonantly enhanced p -wave Fermi gases.

DOI: [10.1103/3mxj-ym4b](https://doi.org/10.1103/3mxj-ym4b)

Understanding pair correlations and their role in fermionic many-body phenomena such as superfluidity is of fundamental importance. Ultracold Fermi gases provide a promising platform to tackle this challenge, owing to the tunability of interactions and thermodynamic quantities. To date, the pair correlations have been quantified by Tan's contact [1–3], which measures short-range correlations between interacting particles and has been extensively characterized in s -wave systems through various experimental studies [4–9]. These studies have clarified the dependence of the contact on the normalized scattering length $1/(k_F a)$ and reduced temperature T/T_F , where k_F and T_F denote Fermi wave number and temperature, respectively. At unitarity $a \rightarrow \infty$, the system exhibits universal behavior solely as a function of T/T_F [10–14] in both harmonically trapped and homogeneous systems, a hallmark of s -wave universal thermodynamics. Precise universal thermodynamic measurements at unitarity [15–18], together with temperature-dependent contact studies, have considerably deepened our understanding of the equation of state for unitary s -wave Fermi gases.

Similarly, for systems with p -wave ($\ell = 1$) interactions, Tan's p -wave contacts have been theoretically formulated [19–29] and measured experimentally in ^{40}K [30–32] and, more recently, in ^6Li [33]. However, no experiments have yet explored their temperature dependence or performed thermodynamic measurements. The temperature

dependence of p -wave contacts is expected to be more complex than in s -wave Fermi gases due to the contribution of the effective range. This follows from the effective range expansion of the p -wave scattering phase shift, $k \cot \delta_p = -1/(vk^2) - 1/R_e$, where v and R_e are the scattering volume and effective range, respectively [34,35]. Near resonance $v \rightarrow \infty$, v dependence vanishes; however, the effective range R_e dominantly determines the phase shift. Thus the contacts and thermodynamics cannot be treated universally, as pointed out in Refs. [22,33].

Another complication in the temperature dependence arises from the dipolar splitting of the p -wave interaction into $m = 0$ and $m = \pm 1$ scattering channels [36,37]. Especially in ^6Li , where the Fermi energy $E_F = k_B T_F$ (k_B : Boltzmann constant) is comparable to the splitting energy δE , the ratio $\delta E/E_F$ serves as a key parameter characterizing the system's anisotropy, as noted in early theoretical studies of the p -wave superfluid phase diagram [38–41] and of p -wave polarons [42]. To capture the temperature dependence of p -wave contacts, it is essential to analyze the contributions from each m channel independently. Besides interaction parameters, a multiple interplay among $k_B T$, E_F , and δE characterizes the p -wave contacts. Thus, once the E_F dependence is established, the $k_B T/E_F$ dependence can be examined—our main objective in this work.

In this Letter, we experimentally determine how the p -wave contacts $C_{v,m}$ in a harmonic trap depend on the Fermi temperature T_F and the reduced temperature T/T_F . By extracting the contacts for the individual m components from the number and dissociation distributions of closed-channel molecules near the p -wave Feshbach resonance (pFBR), we first find that the $m = 0, +1, -1$ components of $C_{v,m}$ exhibit similar behavior as functions of the normalized dimer energy. Moreover, near resonance, the

*Contact author: mukaiyama@phys.sci.isct.ac.jp

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI.

contact increases linearly with the normalized effective range $k_F R_e$ at fixed T/T_F . Finally, we show that the T/T_F dependence reveals opposite temperature trends between the relatively near-resonant and detuned regimes, in agreement with estimates based on the second-order virial expansion. The quantitative comparison with these estimates establishes valuable benchmarks for describing the thermodynamics of p -wave gases and provides insights into their equation of state.

The p -wave contact $C_{v,m}$ can be measured from the ratio of closed-channel molecules to the total number of atoms when the system is held near a pFBR. The molecular fraction $f_{c,m}$ in the $m = 0, +1, -1$ channels is linked to $C_{v,m}$ [22,30]:

$$f_{c,m} = \ell_{c,m}^{-1} C_{v,m} / 2N, \quad (1)$$

where $\ell_{c,m} = M \delta \mu_m v_m^{\text{bg}} \Delta B_m / \hbar^2$. Here, M is the atomic mass, $\delta \mu_m$ is the difference in magnetic moment between the closed- and open-channel states, and $v_m^{\text{bg}} \Delta B_m$ is the product of the background scattering volume and the resonance width. Because the $m = \pm 1$ channels are degenerate and unresolved, we treat them as equally occupied when extracting $C_{v,m}$. In this work, we also exclude the second contact that originates from changes in the effective range [20].

We first prepare spin-polarized ^6Li atoms in the hyperfine state $|2\rangle \equiv |F=1/2, m_F=-1/2\rangle$ in a 1064 nm cigar-shaped optical dipole trap (ODT). The magnetic field is set near 159.1 G, where the pFBR for $|1\rangle-|1\rangle$ ($|1\rangle \equiv |F=1/2, m_F=1/2\rangle$) occurs, with a field uncertainty of 2.5 mG. The atom number is fixed at $N = 1.1(2) \times 10^6$, while T/T_F and T_F are varied independently by adjusting the trap depth. With trapping frequencies of $(\omega_x, \omega_y, \omega_z) = 2\pi \times (660, 9.5, 810)$ Hz at the lowest ODT intensity, we reach $T/T_F = 0.10(2)$ and $T_F = (6N\hbar^3 \omega_x \omega_y \omega_z)^{1/3} / k_B = 1.5(1) \mu\text{K}$ determined by atomic density.

To measure the contacts, we quench the p -wave interaction using a 6.5 μs radio-frequency (rf) π -pulse transferring all atoms from $|2\rangle$ to $|1\rangle$. After a hold time of 0.2–1.5 ms, allowing saturation of the closed-channel fraction, an identical rf pulse transfers only free atoms back to $|2\rangle$; molecules remain in $|1\rangle$ because of negligible wave function overlap. The molecular fraction $f_{c,m}$ is then obtained via state-selective imaging. Because molecules are invisible to imaging light, they are dissociated into atoms by a 1.35 G magnetic-field ramp-up within 20 μs . This rapid ramp-up not only enables molecule detection but also projects their angular momentum distribution onto the absorption image [37,43–46]. After 1 ms time of flight, the dissociated molecules in $|1\rangle$ are imaged at 160.5 G along the y axis, projecting the y -integrated distribution onto the x - z plane with z aligned to the magnetic field.

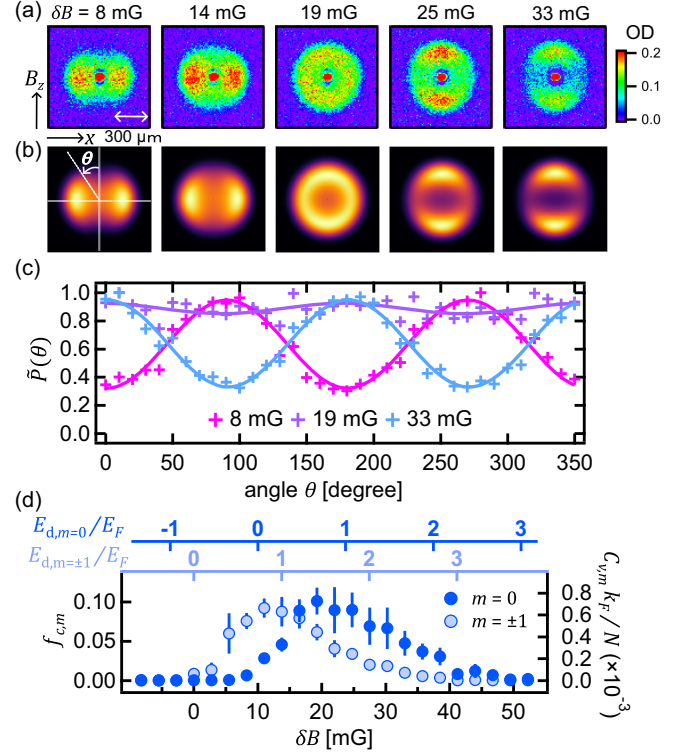


FIG. 1. Observation of the dipolar splitting. (a),(b) Momentum distributions of dissociated p -wave closed-channel molecules in the x - z plane at various magnetic fields: (a) experimental data averaged over 3–6 shots and (b) numerical calculation. The central peaks in (a) originate from open-channel atoms transferred to state $|2\rangle$. (c) Normalized angular probability distribution $\bar{P}(\theta)$ of dissociated molecules. Fits to Eq. (2) (solid lines) extract the anisotropy parameter 2γ , which is 0.99(2), 0.64(1), and 0.35(1) for $\delta B = 8, 19$, and 33 mG, respectively. (d) Fraction of closed-channel molecules $f_{c,m}$ and normalized contacts $C_{v,m} k_F / N$ as functions of magnetic field detuning δB or normalized dimer energy for each m component. Error bars represent the standard deviations of three repetitions.

The number of remaining $|2\rangle$ atoms is measured at 0 G with an arbitrary time of flight.

Figure 1(a) shows the momentum distribution of dissociated closed-channel molecules at various magnetic-field detunings $\delta B = B - B_0$, where B_0 is the resonance field for the $m = \pm 1$ channel. The anisotropy originating from dipolar splitting is evident: as δB increases, the distribution evolves from x to z oriented, passing through an isotropic, center-hollowed shape. This evolution reflects the crossover from the $m = \pm 1$ dominated regime to the $m = 0$ regime, with both contributing in the intermediate region.

To quantify the observed anisotropy, we plot the normalized number of molecules as a function of the polar angle θ in the x - z plane measured from the z axis (0°), as shown in Fig. 1(c). The anisotropy is calibrated as the fraction of $m = +1$ molecules: $\gamma = N_{c,m=+1} / (N_{c,m=0} + N_{c,m=+1} + N_{c,m=-1})$. Because dissociated molecules follow the spherical harmonic distribution for $\ell = 1$,

integrating over the y axis yields the $x - z$ distribution (see Supplemental Material [47]),

$$\tilde{P}(\theta) = C(\gamma) \left[\gamma + (1 - 3\gamma) \times \frac{2}{3} \cos^2 \theta \right], \quad (2)$$

where $C(\gamma)$ is a normalization factor. Fitting Fig. 1(c) with Eq. (2) yields γ . Using this value, we numerically reproduce the dissociated molecular distributions for each image in Fig. 1(b) [47]. The agreement between experiment and simulation confirms the analysis and the resolved dipolar splitting.

While the narrow dipolar splitting of the pFBR in ^6Li has previously been observed in atom-loss spectroscopy in a $\delta E/E_F \gg 1$ regime by reducing the Fermi temperature $T_F = E_F/k_B$ significantly to the order of 100 nK [43,48], our measurements combined with a fast B sweep can resolve the splitting even at the $\delta E/E_F \leq 1$ regime. Remarkably, even at a high Fermi temperature of 1.5 μK , we see fully polarized $m = 0$ or $m = \pm 1$ molecules at the corresponding magnetic fields.

Considering the population balance γ in each m channel, we extract the molecular fraction $f_{c,m}$ and normalized contact $C_{v,m}k_F/N$ from Eq. (1) as functions of the magnetic field [Fig. 1(d)]. The molecular fraction $f_{c,m} = N_{c,m}/N$ is converted to the contact using $\delta\mu_m = k_B \times 113(7) \mu\text{K}/\text{G}$ [49] and $v_m^{\text{bg}} \Delta B_m = -2.8(3) \times 10^6 a_0^3$ [50] (a_0 : Bohr radius), assuming that both parameters are identical for each m . Up to $N_{c,m}/N \sim 10(2)\%$ of atoms are converted into molecules commonly for the three channels, corresponding to a maximum $C_{v,m}k_F/N = 0.7(1) \times 10^{-3}$ at $T/T_F = 0.1$ with $T_F = 1.5 \mu\text{K}$. Notably, the $m = 0$ peak appears at a higher field by 10(2) mG than $m = \pm 1$, consistent with the result of the atom-loss spectroscopy [48]. To represent this splitting, the top axis in Fig. 1(d) scales the dimer binding energy $E_{d,m} = -\hbar^2 R_e/(Mv_m) = \delta\mu_m \delta B$ [40,51], with $R_e = 11a_0$ [52], normalized by E_F . In this representation, $C_{v,m}(m = 0, +1, -1)$ rises sharply near $E_{d,m}/E_F = 0$ and vanishes at ≈ 2 , consistent to previous work in ^{40}K [30]. Contacts for $E_{d,m}/E_F < 0$ are not obtained in this work because the scattering channel is inaccessible to free atoms.

Next, we study how the variations in the Fermi temperature $T_F = E_F/k_B$ affect the anisotropy and the p -wave contacts while keeping $T/T_F = 0.3(1)$ constant. Varying T_F from 1.5 to 6 μK changes the normalized splitting energy $\delta E/E_F$. As shown in Fig. 2(a), the dipolar splitting visible at low T_F —as indicated by the color variation in the map—progressively disappears at higher T_F . In particular, for $T_F \gtrsim 4.5 \mu\text{K}$ ($\delta E/E_F \lesssim 0.25$), we find that the dissociation distribution becomes nearly isotropic, yielding $2\gamma \approx 2/3$ for all magnetic fields. This is a natural behavior in high T_F regimes, where E_F largely exceeds the splitting energy; thus three channels are resonantly excited simultaneously. Here, we extract the relative contributions of the

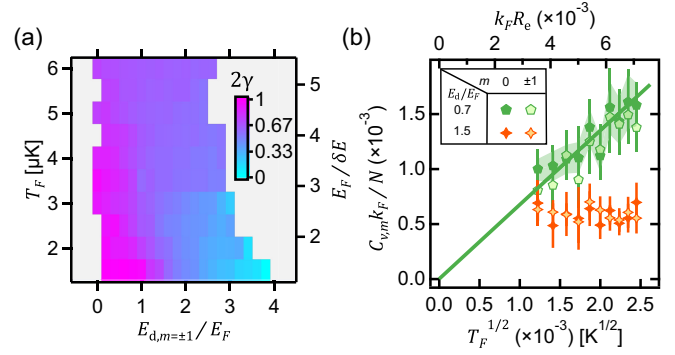


FIG. 2. Fermi temperature dependence. (a) Anisotropy parameter γ versus dimer energy $E_{d,m=\pm 1}/E_F$ at various Fermi temperatures T_F , and the inverse normalized splitting energy $E_F/\delta E$ (right axis). (b) $C_{v,m}$ versus $\sqrt{T_F}$ and the effective range $k_F R_e$ at $E_{d,m}/E_F = 0.7$ (green) and 1.5 (orange). The shaded area indicates a systematic uncertainty arising from atom losses, shown for $m = 0$ and expected to be similar for $m = \pm 1$. The solid line is a guide to the eye.

three channels in this crossover regime, which informs predictions of the p -wave superfluid phase diagram at higher T_F .

The Fermi temperature also affects the p -wave contacts. Figure 2(b) shows $C_{v,m}$ as a function of $\sqrt{T_F}$ near and slightly away from resonance. Near resonance at $E_d/E_F = 0.7$, the contact increases linearly with $\sqrt{T_F}$, and hence with $k_F R_e$, reflecting the dominant contribution of the effective range in which the scattering volume nearly diverges. This $\sqrt{T_F}$ dependence indicates that, unlike the universal behavior of the s -wave contact at unitarity, normalization by $k_F R_e$ is required for universal scaling near resonance, consistent with theoretical predictions [20,24]. By contrast, in the slightly detuned regime at $E_d/E_F = 1.5$, the contact is insensitive to T_F , indicating a breakdown of the near-resonant scaling. Although high-temperature estimates predict a $(k_F R_e)^2$ dependence and hence linear T_F scaling at large detuning [20], such behavior is absent at $E_d/E_F = 1.5$, likely because this detuning is not far enough from resonance. In this work, the large-detuning scaling could not be identified due to the limited signal-to-noise ratio at large detuning.

Reflecting the T_F dependence of the p -wave contacts, we examine their variation with reduced temperature T/T_F . Figure 3 shows $C_{v,m}$ versus dimer energy for various T/T_F at a fixed T_F . We find that, as the system enters deeper into the degenerate regime, the contact in the $E_d/E_F \lesssim 2$ regime increases, whereas the contact in the $E_d/E_F \gtrsim 2$ regime vanishes, yielding a sharper distribution. A distinct feature is observed at high temperatures; the peak decreases and the contact spreads toward larger detuning, so the two regimes exhibit opposite temperature dependences.

This feature can be quantitatively described using the quantum virial expansion of the thermodynamic potential

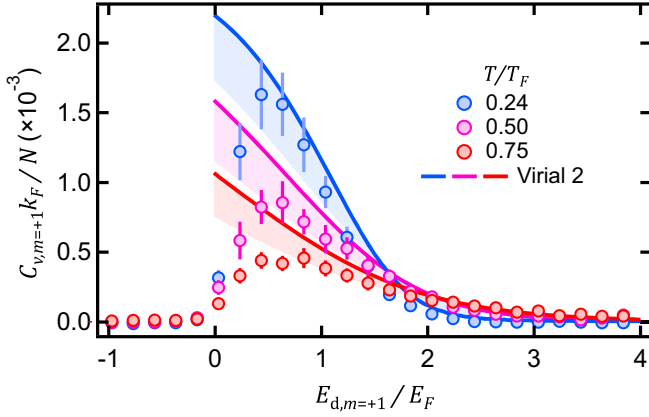


FIG. 3. p -wave contact $C_{v,m=+1}$ versus dimer energy $E_{d,m=+1}/E_F$ for various T/T_F at fixed $T_F = 6.8 \mu\text{K}$ ($k_F R = 7.6 \times 10^{-3}$). Solid lines show second-virial estimates. Results for $m = 0, -1$ are omitted as they agree with $m = +1$. The shaded bands show systematic uncertainties.

in a harmonic trap [53–55]. Here, we use the mean trapping frequency $\omega = (\omega_x \omega_y \omega_z)^{1/3}$. At high temperatures, the thermodynamic potential of an interacting Fermi gas can be expanded as

$$\Omega = \Omega^{(1)} - k_B T Q_1 [\Delta b_{2,p} z^2 + \Delta b_{3,p} z^3 + \dots], \quad (3)$$

where $\Omega^{(1)}$ is the noninteracting thermodynamic potential, $\Delta b_{n,p}$ is the n th p -wave virial coefficient, and $z \equiv \exp(\mu/k_B T)$ (μ : chemical potential) is the fugacity. $Q_1 = (k_B T/\hbar\omega)^3$ is the single-particle partition function for spinless fermions in a harmonic trap. Applying the adiabatic sweep theorem, $\delta\Omega/\delta(v^{-1}) = -\hbar^2 C_v/2M$ ($C_v \equiv \Sigma_m C_{v,m}$) [20,21], and truncating the expansion at second order, we obtain

$$\frac{C_v k_F}{N} = 24\pi (k_F R) \left(\frac{T}{T_F}\right)^3 c_{2,p} z^2, \quad (4)$$

where $c_{2,p} \equiv \partial \Delta b_{2,p} / \partial (R \lambda^2 / v)$ and $\lambda \equiv \sqrt{2\pi \hbar^2 / M k_B T}$. The fugacity z is determined from the number equation

$$\tilde{\rho} = \tilde{\rho}^{(1)}(z) + 2\Delta b_{2,p} z^2, \quad (5)$$

where $\tilde{\rho} \equiv (T_F/T)^3/6$ and $\tilde{\rho}^{(1)}(z) \equiv (1/2) \int_0^\infty dt t^2 / (1 + z^{-1} e^t)$ is the noninteracting density in a harmonic trap.

To evaluate $c_{2,p}$, we use the expression for $\Delta b_{2,p}$ in an isotropic harmonic trap [56,57]:

$$\Delta b_{2,p} = \frac{2\ell + 1}{2} \sum_n [e^{-E_{\text{rel},n}/k_B T} - e^{-E_{\text{rel},n}^{(1)}/k_B T}], \quad (6)$$

where $E_{\text{rel},n} = (2\nu_n + 5/2)\hbar\omega$ and ν_n satisfies $\Gamma(-\nu_n)/\Gamma(-\nu_n - 3/2) = -d^3/8v - (d/4R)(2\nu_n + 3/2)$ [56,58],

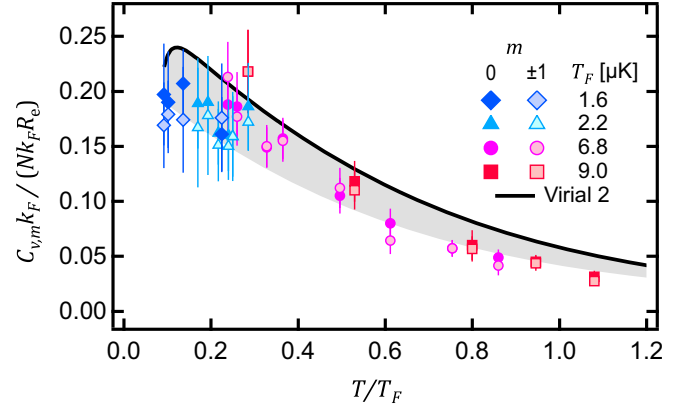


FIG. 4. Peak contact values for $m = 0, +1, -1$ at $E_d/E_F = 0.7$, normalized by $k_F R_e$ at various T_F , plotted as a function of T/T_F . Temperatures are measured before the interaction quench. The black dashed line shows the second-virial estimate with a shaded systematic uncertainty band.

with Γ being the gamma function and $d = \sqrt{2\hbar/M\omega}$ being the characteristic trap length. $E_{\text{rel},n}^{(1)} = (2n + 5/2)\hbar\omega$ ($n = 0, 1, 2, \dots$) is the noninteracting relative energy. At $v^{-1} = 0$, we find $c_{2,p,\infty} = 3/(4\pi)$, and thus Eq. (4) gives $C_v k_F/N = 1/2 (k_F R_e) (T/T_F)^{-3}$ in the high-temperature limit $T \gg T_F$. This contrasts with the homogeneous case, where $C_v \propto (T/T_F)^{-3/2}$ [20], reflecting the faster decay at high T in a harmonic trap due to the peak density reduction [53,54]. We compute $\Delta b_{2,p}$ and its derivative $c_{2,p}$ for any interaction strength to illustrate predicted contact in Fig. 3 for the experimentally measured temperatures, assuming that the three $m = 0, +1, -1$ components of $C_{v,m}$ are equally populated.

This estimation captures the reversal of the temperature dependence in the two detuning regimes separated by the boundary $E_d \simeq 2E_F$ and agrees well with the data for $E_d/E_F \gtrsim 2$. As the system approaches resonance, deviations emerge and become large for $E_d/E_F \lesssim 0.5$. A primary cause is strong atom losses and the associated heating; up to 40% loss and $\Delta(T/T_F) \sim +0.2$ are observed near resonance. The resulting systematic uncertainty in T/T_F , shown as shaded bands, accounts for the deviation down to $E_d/E_F \sim 0.5$. Closer to resonance, strong losses prevent thermalization, making equilibrium-based virial estimates unreliable. A similar reduction of the contact in a loss-dominated regime was reported in Ref. [30], highlighting the need for theoretical models that incorporate nonequilibrium dynamics for more understanding. Another likely source of deviation is the absence of higher-order virial coefficients for three- and four-body processes, which are expected to contribute especially near the resonance.

Figure 4 shows the contact values measured at $E_d/E_F = 0.7$, where the maximum contact is observed over a wide temperature range. Normalizing by $k_F R_e$, corresponding to measuring the slope in Fig. 2(b), collapses

the data for various T_F onto a single curve, indicating universal temperature dependence. When normalized, the increase in the contact toward lower temperatures is evident. For $T/T_F \lesssim 0.3$, the observed contacts appear to saturate or slightly decrease upon further cooling, in qualitative agreement with the prediction by the Nozières–Schmitt-Rink theory [24,29]. Furthermore, we find that the second-virial coefficient seems to reproduce the data even in the deeply degenerate regime $T/T_F < 0.5$, as also seen in Fig. 3, although its applicability is limited to $T/T_F \sim 0.3$ – 0.5 in the s -wave case [12,53,54]. The comparison in this temperature range should be interpreted with caution as the virial expansion becomes less reliable at large fugacity and atom losses can alter the thermodynamics. Our data thus provide a useful benchmark for assessing how far toward the low-temperature side the virial expansion remains applicable in p -wave Fermi gases.

The measured peak contact can be compared with earlier work. Reference [33] reported $C_v/(NR_e) = 0.7$ in ^6Li at $T/T_F \sim 0.3$ as the sum over unresolved m channels, consistent with our summed value of 0.6(1) and strongly validating the use of closed-channel molecule numbers. Reference [30] found $C_{v,m}/(NR_e) \approx 1$ in ^{40}K at $T/T_F \sim 0.2$ for both $m = 0$ and the sum of $m = \pm 1$, whereas we observe 0.2 and 0.4, respectively. The differences remain a missing piece toward a universal understanding.

In conclusion, we have shown that p -wave contacts $C_{v,m}$ are governed by both T_F and T/T_F due to the presence of dipolar splitting and the effective range contributions, in contrast to s -wave Fermi gases. Nevertheless, similar to the s -wave case, the virial expansion provides a powerful framework for estimating thermodynamics. Our results establish a key means for characterizing the many-body properties and thermodynamic behavior of Fermi gases with p -wave interactions.

Acknowledgments—We acknowledge Shi-Guo Peng for fruitful discussion on the second-virial coefficient. This work is supported by JSPS KAKENHI Grant No. JP24K00553. K.N. is supported by the Tsubame Scholarship of Institute of Science Tokyo.

Data availability—The data that support the findings of this article are openly available [59].

[1] S. Tan, *Ann. Phys. (Amsterdam)* **323**, 2952 (2008).
 [2] S. Tan, *Ann. Phys. (Amsterdam)* **323**, 2971 (2008).
 [3] S. Tan, *Ann. Phys. (Amsterdam)* **323**, 2987 (2008).
 [4] J. T. Stewart, J. P. Gaebler, T. E. Drake, and D. S. Jin, *Phys. Rev. Lett.* **104**, 235301 (2010).
 [5] G. B. Partridge, K. E. Strecker, R. I. Kamar, M. W. Jack, and R. G. Hulet, *Phys. Rev. Lett.* **95**, 020404 (2005).
 [6] E. D. Kuhnle, H. Hu, X.-J. Liu, P. Dyke, M. Mark, P. D. Drummond, P. Hannaford, and C. J. Vale, *Phys. Rev. Lett.* **105**, 070402 (2010).

[7] S. Laurent, M. Pierce, M. Delehaye, T. Yefsah, F. Chevy, and C. Salomon, *Phys. Rev. Lett.* **118**, 103403 (2017).
 [8] Y. Sagi, T. E. Drake, R. Paudel, R. Chapurin, and D. S. Jin, *Phys. Rev. Lett.* **114**, 075301 (2015).
 [9] C. Shkedrov, Y. Florshaim, G. Ness, A. Gandman, and Y. Sagi, *Phys. Rev. Lett.* **121**, 093402 (2018).
 [10] Y. Sagi, T. E. Drake, R. Paudel, and D. S. Jin, *Phys. Rev. Lett.* **109**, 220402 (2012).
 [11] S. Hoinka, M. Lingham, K. Fenech, H. Hu, C. J. Vale, J. E. Drut, and S. Gandolfi, *Phys. Rev. Lett.* **110**, 055305 (2013).
 [12] E. D. Kuhnle, S. Hoinka, P. Dyke, H. Hu, P. Hannaford, and C. J. Vale, *Phys. Rev. Lett.* **106**, 170402 (2011).
 [13] B. Mukherjee, P. B. Patel, Z. Yan, R. J. Fletcher, J. Struck, and M. W. Zwierlein, *Phys. Rev. Lett.* **122**, 203402 (2019).
 [14] K. G. S. Xie, C. J. Dale, K. P. Grehan, M. F. Wang, T. Enss, P. S. Julienne, Z. Yu, and J. H. Thywissen, *arXiv:2506.13707*.
 [15] S. Nascimbène, N. Navon, K. J. Jiang, F. Chevy, and C. Salomon, *Nature (London)* **463**, 1057 (2010).
 [16] N. Navon, S. Nascimbène, F. Chevy, and C. Salomon, *Science* **328**, 729 (2010).
 [17] M. Horikoshi, S. Nakajima, M. Ueda, and T. Mukaiyama, *Science* **327**, 442 (2010).
 [18] M. J. H. Ku, A. T. Sommer, L. W. Cheuk, and M. W. Zwierlein, *Science* **335**, 563 (2012).
 [19] M. He, S. Zhang, H. M. Chan, and Q. Zhou, *Phys. Rev. Lett.* **116**, 045301 (2016).
 [20] Z. Yu, J. H. Thywissen, and S. Zhang, *Phys. Rev. Lett.* **115**, 135304 (2015).
 [21] Z. Yu, J. H. Thywissen, and S. Zhang, *Phys. Rev. Lett.* **117**, 019901(E) (2016).
 [22] S. M. Yoshida and M. Ueda, *Phys. Rev. Lett.* **115**, 135303 (2015).
 [23] S.-G. Peng, X.-J. Liu, and H. Hu, *Phys. Rev. A* **94**, 063651 (2016).
 [24] D. Inotani and Y. Ohashi, *Phys. Rev. A* **98**, 023603 (2018).
 [25] F. Qin, X. Cui, and W. Yi, *Phys. Rev. A* **94**, 063616 (2016).
 [26] Y.-C. Zhang and S. Zhang, *Phys. Rev. A* **95**, 023603 (2017).
 [27] X. Yin, X.-W. Guan, Y. Zhang, H. Su, and S. Zhang, *Phys. Rev. A* **98**, 023605 (2018).
 [28] J. Maki, C. J. Dale, J. H. Thywissen, and S. Zhang, *Phys. Rev. A* **110**, 053314 (2024).
 [29] J. Yao and S. Zhang, *Phys. Rev. A* **97**, 043612 (2018).
 [30] C. Luciuk, S. Trotzky, S. Smale, Z. Yu, S. Zhang, and J. H. Thywissen, *Nat. Phys.* **12**, 599 (2016).
 [31] K. G. Jackson, C. J. Dale, J. Maki, K. G. S. Xie, B. A. Olsen, D. J. M. Ahmed-Braun, S. Zhang, and J. H. Thywissen, *Phys. Rev. X* **13**, 021013 (2023).
 [32] C. J. Dale, K. G. S. Xie, K. Pond Grehan, S. Zhang, J. Maki, and J. H. Thywissen, *Phys. Rev. A* **110**, L051302 (2024).
 [33] K. Nagase, S. Oshima, H. Takahashi, and T. Mukaiyama, *Phys. Rev. A* **111**, 013314 (2025).
 [34] J. R. Taylor, *Scattering Theory: The Quantum Theory of Nonrelativistic Collisions* (Wiley, New York, 1972).
 [35] L. D. Landau and E. M. Lifshitz, *Quantum Mechanics: Non-Relativistic Theory*, Course of Theoretical Physics Vol. 3 (Butterworth-Heinemann, London, 1999).
 [36] C. Ticknor, C. A. Regal, D. S. Jin, and J. L. Bohn, *Phys. Rev. A* **69**, 042712 (2004).

-
- [37] J. P. Gaebler, J. T. Stewart, J. L. Bohn, and D. S. Jin, *Phys. Rev. Lett.* **98**, 200403 (2007).
 - [38] M. Iskin and C. A. R. Sá de Melo, *Phys. Rev. Lett.* **96**, 040402 (2006).
 - [39] C.-H. Cheng and S.-K. Yip, *Phys. Rev. Lett.* **95**, 070404 (2005).
 - [40] V. Gurarie and L. Radzihovsky, *Ann. Phys. (Amsterdam)* **322**, 2 (2007).
 - [41] M. Y. Kagan and D. V. Efremov, *J. Exp. Theor. Phys.* **110**, 426 (2010).
 - [42] J. Levinsen, P. Massignan, F. Chevy, and C. Lobo, *Phys. Rev. Lett.* **109**, 075302 (2012).
 - [43] S. Peng, S. Peng, L. Ren, S. Liu, B. Liu, J. Li, and L. Luo, *arXiv:2505.22409*.
 - [44] M. Waseem, Z. Zhang, J. Yoshida, K. Hattori, T. Saito, and T. Mukaiyama, *J. Phys. B* **49**, 204001 (2016).
 - [45] T. Volz, S. Dürr, N. Syassen, G. Rempe, E. van Kempen, and S. Kokkelmans, *Phys. Rev. A* **72**, 010704(R) (2005).
 - [46] N. R. Thomas, N. Kjærgaard, P. S. Julienne, and A. C. Wilson, *Phys. Rev. Lett.* **93**, 173201 (2004).
 - [47] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/3mxj-ym4b> for more details.
 - [48] M. Gerken, B. Tran, S. Häfner, E. Tiemann, B. Zhu, and M. Weidemüller, *Phys. Rev. A* **100**, 050701(R) (2019).
 - [49] J. Fuchs, C. Ticknor, P. Dyke, G. Veeravalli, E. Kuhnle, W. Rowlands, P. Hannaford, and C. J. Vale, *Phys. Rev. A* **77**, 053616 (2008).
 - [50] T. Nakasuji, J. Yoshida, and T. Mukaiyama, *Phys. Rev. A* **88**, 012710 (2013).
 - [51] M. Waseem, T. Saito, J. Yoshida, and T. Mukaiyama, *Phys. Rev. A* **96**, 062704 (2017).
 - [52] M. Waseem, J. Yoshida, T. Saito, and T. Mukaiyama, *Phys. Rev. A* **98**, 020702(R) (2018).
 - [53] X.-J. Liu, H. Hu, and P. D. Drummond, *Phys. Rev. Lett.* **102**, 160401 (2009).
 - [54] H. Hu, X.-J. Liu, and P. D. Drummond, *New J. Phys.* **13**, 035007 (2011).
 - [55] X.-J. Liu, *Phys. Rep.* **524**, 37 (2013).
 - [56] S.-G. Peng, S.-Q. Li, P. D. Drummond, and X.-J. Liu, *Phys. Rev. A* **83**, 063618 (2011).
 - [57] C.-X. Zhang, S.-G. Peng, and K. Jiang, *Phys. Rev. A* **98**, 043619 (2018).
 - [58] K. Kanjilal and D. Blume, *Phys. Rev. A* **70**, 042709 (2004).
 - [59] Kenta Nagase, Dataset for “Temperature dependence of p-wave contacts in a harmonically trapped Fermi gas”, Zenodo, v1, [10.5281/zenodo.17901716](https://zenodo.org/record/17901716) (2025).