Constraints on a possible dineutron state from pionless EFT

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Abstract

We investigate the sensitivity of the three-nucleon system to changes in the neutron–neutron scattering length to next-to-leading order in the pionless effective field theory, focusing on the triton–\textsuperscript{3}He binding energy difference and neutron–deuteron elastic scattering. Due to the appearance of an electromagnetic three-body counterterm at this order, the triton–\textsuperscript{3}He binding energy difference remains consistent with the experimental value even for large positive neutron–neutron scattering lengths while the elastic neutron–deuteron scattering phase shifts are insensitive. We conclude that a bound dineutron cannot be excluded to next-to-leading order in pionless EFT.
I. INTRODUCTION

The search for dineutron bound states has a long history in physics. Although early experimental searches were negative \cite{1, 2}, there has been some evidence for the presence of dineutron configurations in the decay of weakly bound nuclei recently. For example, Bokharev and collaborators claim that roughly half of the excited-state decay of $^6\text{He}$ is through the dineutron \cite{3}, Seth and Parker found evidence for the presence dineutrons in the breakup of $^5\text{H}$, $^6\text{H}$, and $^8\text{He}$ \cite{4}, and Spyrou et al. observed dineutron emission in the ground state decay of $^{16}\text{Be}$ \cite{5}, to mention a few. Whether such dineutron configurations correspond to dineutron bound states, however, is unclear.

In free space the dineutron is believed to be unbound by about 100 keV, implying a large negative scattering length of about $-20$ fm. The most precise determination of the neutron-neutron scattering length to date probably comes from the final-state interaction in the $\pi^-d$ radiative capture reaction \cite{6}, leading to the value $a_{\text{n-n}} = -18.63 \pm 0.27$ (expt.) $\pm 0.30$ (th.) fm \cite{7}. However, the final state interaction peak is expected to be insensitive to the sign of $a_{\text{n-n}}$ such that a positive value of roughly equal magnitude would not be excluded \cite{8}. This issue requires further study.

Because it is just barely unbound, only a small change in the nucleon–nucleon interaction is sufficient to create a bound dineutron. In lattice QCD calculations at unphysically large pion masses of order 800 MeV, e.g., the spin-singlet nucleon–nucleon system and thus the dineutron is bound by about 20 MeV \cite{9, 10}. Moreover, a relatively small change in the quark masses, as it is discussed in scenarios for the variation of fundamental constants, might already be enough to stabilize the dineutron. Kneller and McLaughlin found that big bang nucleosynthesis is compatible with dineutron binding energies of up to 2.5 MeV, thus providing surprisingly weak constraints \cite{11}.

In the context of the nuclear few-body problem, Witala and Glöckle raised the possibility that a slightly bound dineutron might solve some open problems in three-body breakup reactions \cite{12}. They changed the neutron–neutron scattering length by multiplying the CD Bonn potential with an overall strength factor ranging from 0.9 to 1.4. One should keep in mind here that this procedure also changes other low-energy scattering parameters besides the scattering length. Witala and Glöckle found that the neutron–deuteron total and differential cross sections do not rule out a bound dineutron. The neutron–neutron final-state interaction configurations measured in Ref. \cite{13}, however, could not simultaneously be reproduced by their rescaled CD Bonn potential. Overall, a dineutron binding energy larger than 100 keV was excluded in their study. These theoretical studies raised interest in new experimental searches for the dineutron. For example at HIGS/TUNL, there is a proposal to measure the neutron–neutron final state interaction in triton photodisintegration \cite{14, 15}.

The pionless effective field theory (EFT) is ideally suited to study the dependence of low-energy nuclear observables on the neutron–neutron scattering length since the latter appears explicitly as a parameter in the theory. The problem of changing other observables as well when rescaling the potential or coupling constants is thus avoided. The theory is applicable for typical momenta below the pion mass and is frequently used to describe low-energy few-nucleon systems (see e.g. Refs. \cite{16, 18} for a reviews and references to earlier work).

Kirscher and Phillips \cite{19} used pionless EFT to compute a model-independent correlation between the difference of the neutron–neutron and (Coulomb-modified) proton–proton scattering lengths and the triton–$^3\text{He}$ binding energy difference. Their calculation was carried
out at leading order (LO) in the pionless EFT but included isospin breaking effects from the physical scattering lengths in different charge channels. They used this correlation to differentiate between different measured values of the neutron–neutron scattering length and extracted a favored value \( a_{n-n} = (-22.9 \pm 4.1) \) fm. They concluded that values outside of this window are not consistent with the experimental difference in binding energies between the triton and \(^3\)He. Thus their analysis excludes a bound dineutron.

Here, we carry out a similar analysis focusing on the triton–\(^3\)He binding energy difference to next-to-leading order (NLO). It was recently shown that a new electromagnetic counterterm enters in the pionless EFT at this order \([20–22]\). Thus the change in the binding energy difference between triton and \(^3\)He can be absorbed by this counterterm. In the next section, we will review some key points of the formalism of pionless EFT to NLO and discuss the integral equations for the triton and \(^3\)He systems. We then present our analysis of three-nucleon observables as well as the naturalness of the counterterm and conclude.

II. FORMALISM

a. Effective Lagrangian. The effective Lagrangian of pionless EFT can be written in the form

\[
\mathcal{L} = N^\dagger \left( iD_0 + \frac{D^2}{2M_N} \right) N - d^\dagger \left[ \sigma_d + \left( iD_0 + \frac{D^2}{4M_N} \right) \right] d - t^A \left[ \sigma_t + \left( iD_0 + \frac{D^2}{4M_N} \right) \right] t^A \\
+ y_d \left( d^\dagger (N^T P_d^i N) + \text{h.c.} \right) + y_t \left( t^A \left( N^T P_t^A N \right) + \text{h.c.} \right) + \mathcal{L}_{\text{photon}} + \mathcal{L}_3, \tag{1}
\]

with the nucleon field \( N \) and two dibaryon fields \( d^i \) (with spin 1 and isospin 0) and \( t^A \) (with spin 0 and isospin 1), corresponding to the deuteron and the spin-singlet isospin-triplet virtual bound state in S-wave nucleon–nucleon scattering. Spin and isospin degrees of freedom are included by treating the field \( N \) as a doublet in both spaces, but for notational convenience we usually suppress the spin and isospin indices of \( N \). The operators \( P_d^i \) and \( P_t^A \) project out the \(^3\)S\(_1\) and \(^1\)S\(_0\) nucleon–nucleon partial waves.

Furthermore, \( \mathcal{L}_{\text{photon}} \) contains the kinetic and gauge fixing terms for the photons, of which we only keep contributions from Coulomb photons. These correspond to a static Coulomb potential between charged particles, but for convenience we introduce Feynman rules for a Coulomb-photon propagator,

\[
\Delta_{\text{Coulomb}}(k) = \frac{i}{k^2 + \Lambda^2}, \tag{2}
\]

which we draw as a wavy line, and factors \((\pm ie \hat{Q})\) for the vertices.

In the spin-doublet S-wave channel where the triton and \(^3\)He reside, a three-body contact interaction is required for renormalization already at leading order in the EFT \([23]\). We write it here in the form given by Ando and Birse \([24]\),

\[
\mathcal{L}_3 = \frac{M_N H(\Lambda)}{3\Lambda^2} N^\dagger \left( y_d^2 (d^i)^\dagger d^i \sigma^i \sigma^i + y_t^2 (t^A)^\dagger t^B \tau^A \tau^B - y_d y_t [(d^i)^\dagger t^A i \sigma^j \tau^A + \text{h.c.}] \right) N, \tag{3}
\]

where \( \sigma^i \) and \( \tau^A \) are Pauli matrices in spin and isospin space, \( \Lambda \) is a momentum cutoff applied in the three-body equations discussed below and \( H(\Lambda) \) a known log-periodic function of the cutoff that depends on a three-body parameter \( \Lambda_\ast \) \([23]\).
b. Scattering equation. The integral equation for the proton-deuteron scattering amplitude in the $^3$He channel is displayed diagrammatically in Fig. 1. It is a three-component quantity that we denote as $T_{\text{full}} = (T_{\text{full}}^d, T_{\text{full}}^{\alpha}, T_{\text{full}}^d)$, where all three components are in general functions of the total energy $E$ as well as of the in- and outgoing momenta $k$ and $p$, i.e., $T = T(E; k, p)$. Everything is projected onto the S-waves here such that there is no angular dependence.

\[
\begin{align*}
T_{\text{full}}^d &= \frac{1}{2\pi^2} \int_0^\Lambda dq \, q^2 A(..., q) B(q, \ldots) \\
T_{\text{full}}^{\alpha} &= \ldots \\
T_{\text{full}}^d &= \ldots
\end{align*}
\]

FIG. 1: Coupled-channel integral equation for the full scattering amplitude in the $^3$He channel. Thin single lines represent nucleons, the double line stands for a deuteron, and the thick lines for dibaryons in the spin-singlet state (with an additional dot to indicate the $p-p$ channel which is treated separately). The diagrams representing the three-nucleon force have been omitted.

Using the formal notation

\[
A \otimes B \equiv \frac{1}{2\pi^2} \int_0^\Lambda dq \, q^2 A(..., q) B(q, \ldots)
\]
Explicit expressions for the kernel functions—denoted by \( K \)—to the existence of a hypothetical bound dineutron, it is more convenient to match the propagators in Eqs. (6) and (7) which are linear in the corresponding effective interactions. The triton channel.

For the channels corresponding to Coulomb-photon exchanges (\( \gamma \) doublet) channel from Eq. (5). As a first step to that end one simply removes all kernel functions in Eq. (5) that correspond to Coulomb-photon exchanges (\( K_{c}^{(d,t)} \), \( K_{box}^{(in/out)} \)). If one furthermore lets \( D_{t}^{pp} \rightarrow D_{t} \), one obtains just the scattering equation for the \( n-d \) doublet amplitude in the isospin limit, which can actually be reduced to a two-channel equation [22].

In this work, however, we want to study the effect of varying the neutron–neutron scattering length, so we rather let \( D_{t}^{pp} \rightarrow D_{t}^{nn} \), where

\[
D_{t,nn}(E; q) = -\frac{4\pi}{M_{N}y_{t}^{2}} \times \frac{1}{-\gamma_{n-n} + \sqrt{3q^{2}/4 - M_{N}E - i\varepsilon}}.
\]

For \( a_{n-n} < 0 \), we simply set \( \gamma_{n-n} \equiv 1/a_{n-n} \). In the regime of positive \( a_{n-n} \), corresponding to the existence of a hypothetical bound dineutron, it is more convenient to match the

\[
\begin{align*}
D_{d,t}(E; q) &= -\frac{4\pi}{M_{N}y_{d,t}^{2}} \times \frac{1}{-\gamma_{d,t} + \sqrt{3q^{2}/4 - M_{N}E - i\varepsilon}} \quad (6) \\
D_{t,pp}(E; q) &= -\frac{4\pi}{M_{N}y_{t}^{2}} \times \frac{1}{-1/aC - \alpha M_{N} \left( \psi(\eta) + \frac{1}{2\eta} - \log(\eta) \right)}.
\end{align*}
\]

Explicit expressions for the kernel functions—\( K_{s}(E; k, p) \), \( K_{c}^{(d,t)}(E; k, p) \), etc.—as well as a more detailed derivation of Eq. (5) can be found in Ref. [22], Secs. III and in particular V.B.2.

\section{Higher-order corrections}

At next-to-leading order there are perturbative corrections to the propagators in Eqs. (9) and (7) which are linear in the corresponding effective ranges (cf. Ref. [20] and, for an expression in the same notation used here, in particular Eq. (3) in Ref. [27]). Our fully perturbative NLO calculation is based on Refs. [28] and [20].

\( d \). The triton channel. We only show the integral equation for the \(^3\text{He} \) (\( p-d \) doublet) channel explicitly. It is straightforward to obtain the integral equation for the \(^3\text{H} \) (\( n-d \) doublet) channel from Eq. (5). As a first step to that end one simply removes all kernel functions in Eq. (5) that correspond to Coulomb-photon exchanges (\( K_{c}^{(d,t)} \), \( K_{box}^{(in/out)} \)). If one furthermore lets \( D_{t}^{pp} \rightarrow D_{t} \), one obtains just the scattering equation for the \( n-d \) doublet amplitude in the isospin limit, which can actually be reduced to a two-channel equation [22].

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\[
D_{t,nn}(E; q) = -\frac{4\pi}{M_{N}y_{t}^{2}} \times \frac{1}{-\gamma_{n-n} + \sqrt{3q^{2}/4 - M_{N}E - i\varepsilon}}.
\]
propagator to the effective range expansion around the dineutron pole. Accordingly, we set
\[ \gamma_{n-n} \equiv \frac{1}{a_{n-n}} \left( 1 + \frac{\rho_{n-n}}{2a_{n-n}} \right) \text{ for } a_{n-n} > 0, \] (9)
where for simplicity we assume \( \rho_{n-n} = \rho_t = 2.73 \text{ fm} \). At leading order, this only corresponds to a constant offset and does not otherwise affect the result. The effect of varying \( \rho_{n-n} \) away from the isospin-symmetric case at NLO will be discussed below.

III. RESULTS

From Eq. (5) and its analog for the neutron–deuteron case one can extract both scattering information—for example the \( n-d \) doublet scattering length which we use as physical input to fix the three-nucleon force \( H(\Lambda) \)—and bound state properties. To extract the binding energies of the triton and \( ^3\text{He} \), we look for poles (as a function of the energy) in the corresponding scattering amplitudes at negative energies. In practice, this can simply be done by studying the homogeneous versions of Eq. (5) and its analog for the triton channel.

A. Leading order

In the left panel of Fig. 2 we show the triton–\(^3\text{He} \) binding energy difference
\[ \Delta E_3 = E_B(^3\text{H}) - E_B(^3\text{He}) \] (10)
as a function of the \( n-n \) scattering length \( a_{n-n} \), both for negative and for positive values of the latter quantity.\(^1\)

In the regime of negative \( a_{n-n} \) our results agree nicely with the findings of Kirscher and Phillips [19], who calculated the binding-energy difference using the resonating group method (RGM). Those authors did not explore the possibility of a (large) positive \( n-n \) scattering length. Consequently, only the negative arm of the pole at \( a_{n-n} = 0 \text{ fm} \) was found in Ref. [19], and from that one might naively think that positive values of \( a_{n-n} \) are clearly excluded. However, the relevant quantity here is not actually \( a_{n-n} \), but rather its inverse. In the right panel of Fig. 2 we show that \( \Delta E_3 \) is indeed a continuous function of \( a_{n-n}^{-1} \) around \( a_{n-n}^{-1} \approx 0 \).

This raises the question of how well one can determine \( a_{n-n} \) from a leading-order pionless EFT calculation. In the left panel of Fig. 2 we show an error band that was generated by varying the (essentially cutoff-converged) \( \Lambda = 1800 \text{ MeV} \) curve within \( \pm 30\% \), corresponding roughly to the estimated size of an NLO correction in pionless EFT. From that, a positive value of \( a_{n-n} \) is just barely excluded, and by making a slightly more conservative estimate one would find that such a case can be marginally consistent with the physical binding-energy difference.

\(^1\) We note that the region of small scattering lengths, \( a_{n-n} \approx 0 \text{ fm} \), should be discarded. This region is clearly excluded by experiment. Moreover, our theory requires a scattering length large compared to the range of the interaction.
FIG. 2: Leading-order triton–$^3$He binding energy difference as a function of the neutron–neutron scattering length (left panel) and as a function of the inverse neutron–neutron scattering length (right panel). The shaded bands were generated by varying the $\Lambda = 1800$ MeV within $\pm 30\%$. The dotted line in the right panel shows the (hypothetical) dineutron energy as a function of $a_{n-n}^{-1}$.

One could argue now that our band might be overestimating the uncertainty since some contributions can be expected to cancel in the difference of the—individually calculated—$^3$H and $^3$He energies. However, it has recently been shown [20] that a new three-body counterterm $H_{0,1}^{(a)}(\Lambda)$ is necessary to renormalize the doublet-channel $p$–$d$ system at next-to-leading order. This might mean that the actual uncertainty is somewhat larger than what one might expect based on the considerations in Ref. [19]. We will investigate this issue in the next subsection.

B. Next-to-leading order

When the new counterterm is fit to reproduce the experimental $^3$He binding energy, one can no longer predict both $E_B(^3\text{H})$ and $E_B(^3\text{He})$. Still, there is a dependence of $\Delta E_3$ on $a_{n-n}$ that comes from the triton binding energy. The result of this calculation is shown in the left panel of Fig. 3.

The shape of the curves is now different compared to the leading-order result, and it looks now as if not only a large range of negative $a_{n-n}$ would be consistent with the experimental $\Delta E_3$, but also some positive values would in fact be allowed. If one looks directly at the prediction for the triton binding energy, as shown in the right panel of Fig. 3, one finds that the results at leading order and next-to-leading order nicely overlap and that the cutoff-dependence is smaller at NLO. To make the figure less cluttered, we show no explicit error bands here. Varying the $n$–$n$ effective range by $\pm 10\%$ around $\rho_{n-n} = \rho_t = 2.73$ fm moves the NLO curves in Fig. 3 up and down by about 0.1-0.15 MeV. The influence of this parameter therefore quite small, but it makes it a little bit harder yet to exclude a bound dineutron on the grounds of pionless effective field theory.

With the new three-body counterterm present one can no longer make a parameter-free prediction for $\Delta E_3$ and next-to-leading order. Unfortunately, fitting $H_{0,1}^{(a)}(\Lambda)$ to the doublet-channel $p$–$d$ scattering length does not restore much predictive power since that quantity is surprisingly poorly known [29].
This makes it interesting to see what can be learned from the $a_{n-n}$-dependence of the new counterterm itself. In particular, one can check if it becomes unnatural for positive values of $a_{n-n}$. In Fig. 4 we answer this question in the negative. The leading behavior is

$H_0^{(\alpha)}(\Lambda) \propto \Lambda$,  \hspace{1cm} (11)

with subleading logarithmic corrections. One clearly sees that the coefficient does not change its order of magnitude if one considers positive values of $a_{n-n}$, neither for natural—$O(m_\pi)$—nor for asymptotically large cutoffs. On this ground we conclude that one cannot rule out the existence of a shallow bound $n-n$ state from pionless EFT at NLO. It might thus be worthwhile to continue investigating that possibility, both theoretically and experimentally.

Finally, we show in Fig. 5 that the $n-d$ scattering phase shifts are quite insensitive to variations of $a_{n-n}$ to (large) positive values. We have only included S-wave phase shifts here since the P-wave result looks very similar and there is even less effect in higher partial waves. Thus there is very little sensitivity in elastic scattering observables to changes from

\[ \begin{align*}
\text{FIG. 3: Left panel: NLO result for the triton–}^3\text{He binding energy difference as a function of the inverse neutron–neutron scattering length. Right panel: NLO result for the triton binding energy as a function of the inverse neutron–neutron scattering length.} \\
\text{FIG. 4: Electromagnetic counterterm } H_0^{(\alpha)}(\Lambda) \text{ as a function of the cutoff } \Lambda \text{ for different neutron-neutron scattering lengths.} \\
\end{align*} \]
FIG. 5: S-wave $n$–$d$ doublet channel scattering phase shifts as functions of the center-of-mass momentum $k$ for a cutoff $\Lambda = 1800$ MeV and several values of the $n$–$n$ scattering length (see plot legend). The shaded bands were generated by letting the $a_{n-n}$ curve vary within ±30% (LO) and ±10% (NLO). The crosses are the results from the AV18+UR potential-model calculation reported in Ref. [30].

$a_{n-n}^{-1}$ from small negative to small positive values which is in agreement with the findings of Witala and Glöckle [12].

IV. CONCLUSION

We conclude that a bound dineutron cannot be excluded based on the triton–$^3$He binding energy difference and elastic scattering results in pionless EFT at NLO. This result provides support for planned dineutron searches by measuring the neutron–neutron final-state interaction in triton photodisintegration [14, 15]. Even if no bound dineutron is found such experiments will be useful to settle the controversy about the value of the neutron–neutron scattering length (cf. Ref. [19]). We note in passing that the existence of a bound dineutron would also provide a way to understand the recent data by the HypHI collaboration suggesting a bound $nn\Lambda$ system [31]. Hiyama and collaborators showed that such a bound state could be accommodated if a bound dineutron state existed [32]. The resulting shift in the triton–$^3$He binding energy difference can be absorbed by a naturally-sized NLO three-body force $H_{0,1}^{(n)}$ as demonstrated above. We note, however, that a bound dineutron would be difficult to accommodate in standard approaches to charge-symmetry breaking in the two-nucleon system [33, 34]. This issue requires further theoretical study.

It would also be valuable to extend the calculation to $N^2$LO and to investigate whether the neutron–neutron final-state interaction configurations measured in Ref. [13] can simultaneously be reproduced with a positive scattering length. The latter would require a calculation of deuteron breakup reactions and is beyond the scope of this work.
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