Chiral three-nucleon forces and the evolution of correlations along the oxygen isotopic chain

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Background: Three-nucleon forces (3NFs) have non trivial implications on the evolution of correlations at extreme proton-neutron asymmetries. Recent ab initio calculations show that leading order chiral interactions are crucial to obtain the correct binding energies and neutron driplines along the O, N, and F chains, Ref. [A. Cipollone, C. Barbieri, P. Navrátil, Phys. Rev. Lett 111, 062501 (2013)].

Purpose: Here, we discuss the impact of 3NFs along the oxygen chain for other quantities of interest, such has the spectral distribution for attachment and removal of a nucleon, spectroscopic factors and matter radii. The objective is to better delineate the general effects of 3NFs on nuclear correlations.

Methods: We employ self-consistent Green’s function (SCGF) theory which allows a comprehensive calculation of the single particle spectral function. For the closed subshell isotopes, 14O, 16O, 22O, 24O and 28O, we perform calculations with the Dyson-ADC(3) method which is the state of the art for both nuclear physic and quantum chemistry applications. The remaining open shell isotopes are investigted using the newly developed Gorkov-SCGF formalism up to second order.

Results: We produce plots for the full-fledged spectral distributions. The spectroscopic factors for the dominant quasiparticle peaks are found to depend very little on the leading order (NNLO) chiral 3NFs. Preexisting 3NFs have small impact on the calculated matter radii, which, however are consistently obtained smaller than experiment. Similarly, single particle spectra tend to be diluted with respect to experiment. This effect might hinder, to some extent, the onset of correlations and screen the quenching of calculated spectroscopic factors. The most important effects of 3NFs is thus the fine tuning of the energies for the dominant quasiparticle states, which govern the shell evolution and the position of driplines.

Conclusions: Although present chiral NNLO 3NFs interactions do reproduce the binding energies correctly in this mass region, the details of the nuclear wave function remain at odd with the experiment showing too small radii and a too dilute single particle spectrum, similar to what already pointed out for larger masses. This suggests a lack of repulsion in the present model of NN+3N interactions which is mildly apparent already for masses in the A=14–28 range.

I. INTRODUCTION

The concept of correlations is fundamental to a deep understanding of nuclear phenomena [1]. These are generally defined as characteristics of the nucleus that cannot be explained in terms of a simple mean field picture (i.e., a wave function of Slater determinant type). These effects are often quantified in terms of the fragmentation of the single particle strength observed for the addition or removal of a nucleon. An intriguing feature is the persistence of dominant quasiparticle peaks near the Fermi surface while broader resonances are found at higher excitations. This is at the origin of the duality between the liquid drop and the shell model behaviour of atomic nuclei.

Historically, several electron scattering studies have provided a wealth of information on nuclear spectral functions (see [2–9] and references therein). This has allowed a rather complete characterisation of correlations for stable nuclei [1]. However, a similar full characterisation for exotic isotopes is still lacking. Recent data from radioactive beam facilities have put in evidence new phenomena such as shell evolution with changing proton-neutron asymmetry [10] and the insurgence of new magic numbers [11–13]. From the theoretical point of view, some of these effects have been explained in terms of properties of the tensor interaction [14] and three-nucleon forces (3NFs) [15]. More recently, it has been shown that 3NFs are crucial for understanding the neutron rich side of the nuclear chart. In particular, to explain the oxygen dripline at 24O [16–19] and neutron rich Ca isotopes [20–23].

Ab-initio calculations of atomic nuclei have advanced dramatically in the medium mass region. Several approaches such as coupled cluster [24, 25], in-medium similarity renormalisation group (IM-SRG) [26] and self-consistent Green’s function (SCGF) [19, 27] theories are now capable of approaching masses up to A≈100 or more. These allow to fully exploit modern chiral interactions with two-nucleon (NN) and 3NFs evolved through SRG techniques [28]. Very recently, open shell nuclei have become accessible through a Gorkov extension of the SCGF formalism [22, 29, 30] and a multi-reference IM-SRG [18, 23]. The SCGF approach has the additional advantage to provide consistent calculations for optical

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potentials and the spectral function for the whole energy spectrum (i.e., both close and far form the Fermi surface). However, all the \textit{ab initio} studies mentioned above have so far focussed on ground state properties, such as total binding energies and two-nucleon separation energies.

In this paper, we consider the Green’s functions of the oxygen isotopes already obtained in Ref. [19] and extend these calculations to the remaining even mass—and open shell—isotopes, using the Gorkov-SCGF approach. We then present first fully microscopic calculations of the evolution of the single particle spectral functions along a full isotopic chain. This gives an overall insight on the evolution of nuclear correlations between two extremes of the nuclear chart. At the same time, it allows to perform a more thorough test of modern chiral interactions and, in particular, we investigate the effects of initial 3NFs at NNLO.

Section II discusses the relevant features of the SCGF formalism, for completeness. It reviews the links of propagators with the spectral function and other quantities of experimental interest. Calculations are done in an \textit{ab initio} fashion and we discuss in some details the choice of the Hamiltonian, the approximations taken and the expected uncertainties, when these can be estimated. This is done in Sec. III. Sec. IV discusses our results for single particle spectra, spectroscopic factors and binding energies. Full three-dimensional plots of spectral functions are shown in the Appendix for completeness and conclusions are drawn in Sec. V.

II. SCGF FORMALISM

Information about the single particle dynamics is fully contained in the one-body Green’s function, or propagator, whose Lehmann representation reads:

\[
g_{\alpha\beta}(\omega) = \sum_n \frac{\langle \Psi_0^A | a_\alpha | \Psi_{n+1}^A \rangle \langle \Psi_{n+1}^A | a_\beta | \Psi_0^A \rangle}{\omega - \varepsilon_n^+ + i\eta} + \sum_k \frac{\langle \Psi_0^A | a_\alpha | \Psi_{k-1}^A \rangle \langle \Psi_{k-1}^A | a_\beta | \Psi_0^A \rangle}{\omega - \varepsilon_k^- - i\eta}.
\]  

In Eq. (1), $|\Psi_0^A\rangle$ represents the ground state of a nucleus and $|\Psi_{n+1}^A\rangle$, $|\Psi_{k-1}^A\rangle$ are the eigenstates of the $(A \pm 1)$-nucleon system. The greek indices $\alpha, \beta, \ldots$, label a complete orthonormal single particle basis, while $\varepsilon_n^+ \equiv (E_{n+1}^A - E_0^A)$ and $\varepsilon_k^- \equiv (E_k^A - E_{k-1}^A)$ are one-nucleon addition and removal energies, respectively. These are generically referred to in the literature as “separation” or “quasiparticle” energies and we will use both names in the following. The transition amplitudes $X_n^\alpha \equiv \langle \Psi_{n+1}^A | a_\alpha | \Psi_0^A \rangle$ and $X_k^\alpha \equiv \langle \Psi_{k-1}^A | a_\alpha | \Psi_0^A \rangle$ give information about the strength of the corresponding particle addition and removal processes.

The one-body Green’s function (1) is completely determined by solving the Dyson equation,

\[
g_{\alpha\beta}(\omega) = g_{\alpha\beta}^0(\omega) + \sum_{\gamma\delta} g_{\alpha\gamma}^0(\omega) \Sigma_{\gamma\delta}^*(\omega) g_{\delta\beta}(\omega),
\]  

where the unperturbed propagator $g_{\alpha\beta}^0(\omega)$ is the initial reference state (usually a mean field or Hartree-Fock state) while $g_{\alpha\beta}(\omega)$ is the correlated propagator. A full knowledge of the self-energy $\Sigma_{\alpha\beta}(\omega)$ yields the exact solution for $g_{\alpha\beta}(\omega)$. However, in practical calculations this has to be approximated and it is expanded in terms of the propagator itself (that is, $\Sigma^* = \Sigma^*[g(\omega)]$). Thus, iterative procedure is required to solve for $\Sigma^*(\omega)$ and Eq. (2) self consistently. The approximation schemes we employ to calculate the self-energy are outlined in the next subsection.

The attractive feature of the SCGF approach is that $g_{\alpha\beta}(\omega)$ describes the one-body dynamics completely. The particle and hole spectral functions are extracted directly from Eq. (1), respectively:

\[
S_{\alpha\beta}^P(\omega) = \sum_n (X_n^\alpha)^* X_n^\beta \delta(\omega - (E_{n+1}^A - E_0^A)),
\]

\[
S_{\alpha\beta}^h(\omega) = \sum_k Y_k^\beta (Y_k^\alpha)^* \delta(\omega - (E_k^A - E_{k-1}^A)).
\]  

Any one-body observable can be calculated via the one-body density matrix $\rho_{\alpha\beta}$, which is obtained from $g_{\alpha\beta}$ as follows:

\[
\rho_{\alpha\beta} \equiv \langle \Psi_0^A | a_\beta | a_\alpha | \Psi_0^A \rangle = \int_{-\infty}^{\infty} S_{\alpha\beta}(\omega) \, d\omega = \sum_k (Y_k^\alpha)^* Y_k^\beta.
\]  

The expectation value of a one-body operator, $\hat{O}^{1B}$, can then be written in terms of the $\rho$ amplitudes as,

\[
\langle \hat{O}^{1B} \rangle = \sum_{\alpha\beta} O_{\alpha\beta}^{1B} \rho_{\beta\alpha} = \sum_k \sum_{\alpha\beta} (Y_k^\alpha)^* O_{\alpha\beta}^{1B} Y_k^\beta.
\]  

Evaluating two- and many-nucleon observables requires the knowledge of many-body propagators. In the following, we do this by approximating the corresponding A-body density matrices with A correlated but non-interacting propagators, Eq. (4). Specifically, we evaluate the matter rms radii as

\[
\langle r^2 \rangle_{\text{rms}} \simeq \frac{1}{A} \left( 1 - \frac{1}{A} \right) \sum_{\alpha\beta} \langle \alpha | \vec{r}^2 | \beta \rangle \rho_{\beta\alpha} + \frac{1}{4A^2} \sum_{\alpha\beta\gamma\delta} \langle \alpha \gamma | \vec{f}_1 \cdot \vec{f}_2 | \beta \delta \rangle \rho_{\beta\alpha} \rho_{\partial\gamma},
\]  

where $\vec{f}_i$ represents the position of particle $i$. The two-body correction term in Eq. (6) arises because the intrinsic radius is calculated with respect to the centre of mass (COM) of the system [31]. The exact one-body propagator, $g_{\alpha\beta}(\omega)$, also allows calculating the total energy by
means of the extended Koltun sum-rule [32]:

\[
E_0^A = \sum_{\alpha\beta} \frac{1}{2} \int_{-\infty}^{\infty} [T_{\alpha\beta} + \omega \delta_{\alpha\beta}] \xi^h_{\beta\alpha} \omega d\omega - \frac{1}{2} \langle W \rangle. \tag{7}
\]

This requires only the additional evaluation of the expectation value of the three-nucleon interaction, \( \langle W \rangle \). Again, we approximate this in terms of non-interacting three-body density matrices:

\[
\langle W \rangle \simeq \frac{1}{6} \sum_{\alpha\beta\gamma\delta\epsilon} W_{\alpha\beta\gamma\delta\epsilon} \rho_{\alpha\gamma} \rho_{\delta\epsilon} \rho_{\epsilon\delta}. \tag{8}
\]

The errors in this approximation have been estimated in Ref. [19] and were found to not exceed the 250 keV on the total binding energy for \(^{16}\)O and \(^{24}\)O.

In all simulations below we subtract the spurious contribution of the kinetic energy of the COM and work with the intrinsic Hamiltonian \( H[A] = H - T_{COM}(A) = U(A) + V(A) + W \), which acquires a dependence on total number of nucleons. The \( U, V \) and \( W \) label one-, two- and three-body interactions. This implies that the particle and hole spectra of the even-odd isotopes are recalculated separately from \( H[A+1] \) and \( H[A-1] \). They are then corrected for the COM motion as follows:

\[
\varepsilon^+_{n,COM} = \varepsilon^+_n [A+1] - E^A_0 [A+1] - E^A_0 [A], \tag{9}
\]

\[
\varepsilon^-_{k,COM} = \varepsilon^-_k [A-1] - E^A_0 [A-1] + E^A_0 [A],
\]

where \( \varepsilon^+_n [A \pm 1] \) and \( E^A_0 [A \pm 1] \) label the poles of \( g_{\alpha\beta}(\omega) \) and the total energies, Eq. (7), calculated from the \( H[A \pm 1] \) Hamiltonian. The overall COM corrections become progressively smaller as \( A \) increases.

### A. Dyson-ADC(3) and second-order Gorkov equations

Calculations with 3N interactions follows the procedure extensively discussed in Ref. [32], which involves defining the following medium dependent one- and two-body interactions:

\[
\bar{U}_{\alpha\beta} = U_{\alpha\beta} + \frac{1}{4} \sum_{\delta\gamma} W_{\alpha\gamma\beta\delta} \rho_{\delta\gamma},
\]

\[
\bar{V}_{\alpha\gamma,\beta\delta} = V_{\alpha\gamma,\beta\delta} + \sum_{\mu\nu} W_{\alpha\mu\gamma,\beta\delta} \rho_{\mu\nu}. \tag{10}
\]

This allows neglecting residual contributions in \( W \) that have been found to be negligible for oxygen isotopes [33, 34]. Hence, we retain only interaction-irreducible diagrams in \( \bar{U} \) and \( \bar{V} \) to the self-energy.

To solve Eq. (2), we express the self-energy as,

\[
\Sigma^*_{\alpha\beta}(\omega) = \Sigma_{\alpha\beta}(\infty) + \sum_j D_{\alpha\beta}^{ij} \left[ \frac{1}{\omega - (K+C)} \right]_{ij} D_{\beta\alpha}^{ji}, \tag{11}
\]

where \( \Sigma_{\alpha\beta}(\infty) \) is the correlated and energy-independent mean field. The whole \( \Sigma_{\alpha\beta}(\omega) \) is an optical potential for elastic scattering of a nucleon off the \( |\Psi_0^{\alpha} \rangle \) ground state, which also describes the fragmentation of the particle and hole spectra [35, 36].

In Eq. (11), the matrix \( D \) couples single particle states to more complex intermediate configurations, while \( K \) and \( C \) are their unperturbed energies and interaction matrices. For the closed subshell isotopes we exploit the third order algebraic diagrammatic construction [ADC(3)] scheme, which is the best compromise between computational efforts and accuracy. This consists in the minimal choice of these matrices that retains all self-energy diagrams up to third order. Generally speaking ADC(\( n \)) defines a hierarchy of truncation schemes of Eq. (11) for increasing order \( n \) that allows systematic improvements of the method [37].

Recently, SCGF theory has been extended to a Nambu-Gorkov formulation that allows addressing truly open shell nuclei [29]. This has opened the possibility to calculate ground state properties and the one-nucleon addition/removal spectra of mid-mass open shell nuclei, in a fully \textit{ab initio} fashion. As in BCS theory, one allows for an explicit breaking of particle-number conservation that is necessary for a proper description of pairing correlations [29, 38]. This implies introducing a grand canonical Hamiltonian \( \Omega = H - \mu_n N - \mu_p \tilde{N} \) and constraining the proton (neutron) chemical potentials \( \mu_p(\mu_n) \) to recover the correct particle number on average: \( A = \langle \Psi_0 | A | \Psi_0 \rangle \), where \(|\Psi_0\rangle \) is the symmetry-broken ground state. A detailed description of the theory can be found in Refs. [29, 30, 39].

In Gorkov theory one is left with a set of normal and anomalous propagators and self-energies with similar Lehman representations to Eqs. (1) and (11). In particular, the normal propagator is

\[
G^{11}_{\alpha\beta}(\omega) = \sum_k \left\{ \frac{U^k_{\alpha\beta} \bar{U}^*_{\alpha\beta}}{\omega - \bar{\omega}_k + i\eta} + \frac{\bar{V}^k_{\alpha\beta} V^*_{\alpha\beta}}{\omega + \bar{\omega}_k - i\eta} \right\}. \tag{12}
\]

where \( U, V \) are the transition amplitudes for reaching the states \(|\Psi_k \rangle \) by adding (removing) a nucleon to (from) \(|\Psi_0\rangle \), and \( \bar{\omega}_k \) are the corresponding separation energies [29].

The Gorkov version of the SCGF approach allows to calculate spectral functions for open shell semi-magic systems. Its present formulation follows the ADC(\( n \)) truncation scheme discussed above but has been implemented only up to the second order. This has allowed successful predictions of trends in binding energies [22]. However, the ADC(2) is known to slightly underestimate binding energies and it is not guaranteed to provide accurate predictions one-nucleon separation energies [40], which are instead possible with a Dyson-ADC(3) calculations. The full extension to Gorkov-ADC(3) formalism is currently underway [41].
III. CALCULATIONS

Calculations have been performed using NN and 3N chiral interactions evolved to low-momentum scales through free space similarity renormalization group (SRG) techniques [28]. The original NN interaction is the next-to-next-to-leading order (N3LO) with a reduced cutoff $\Lambda_{NN} = 500$ MeV/$c$, from Refs. [42, 43]. For the 3N interactions we used the NNLO with a reduced local cutoff of $\Lambda_{3N} = 400$ MeV/$c$ [34, 44]. This includes the two-pion exchange contribution that was originally proposed by Fujita and Miyazawa [45]. Low-energy constants were set at $c_D = -0.2$, $c_E = 0.098$ to reproduce the binding energy of $^4$He. When we perform the SRG transformation of the sole NN-N3LO interaction we already obtain evolved NN+3N interactions. We will refer to this as the “induced” Hamiltonian. Conversely, the “full” Hamiltonian is the one obtained by also evolving the Fujita-Miyazawa type are included in the full Hamiltonian only.

All calculations were performed in a model space of 12 harmonic oscillator (HO) shells [$N_{\text{max}} \equiv \max(2n + l) = 11$], including all NN matrix elements and limiting 3NF ones to configurations with $N_1 + N_2 + N_3 \leq 14$. Changing the oscillator frequency between $\hbar \Omega_{\text{HO}} = 20$ and 24 MeV in Dyson-ADC(3) calculations, we found up to 450 KeV variations in the binding energy of $^{24}$O. Similarly, varying $\lambda_{\text{SRG}}$ in a limited range $1.88 - 2.0$ fm$^{-1}$ did not induce variations of more than 0.5%. From these and other tests we infer a conservative theoretical error of at most 5%, for binding energies obtained with Dyson-ADC(3) [19]. Similar conclusions can be drawn about the prediction of dominant quasiparticle peaks in the single particle spectrum, $\epsilon^+_n$ and $\epsilon^-_n$. Varying both $\hbar \Omega_{\text{HO}} = 20$-24 MeV and $\lambda_{\text{SRG}} = 1.88$-2.0 fm$^{-1}$, we found a maximum variation of 310 KeV for the neutron $1/2^-$ quasihole in $^{24}$O. This corresponds to 2% of its value, $\epsilon^{-}_{1/2^-} = -14.22$ MeV. The largest variation for proton separation energies was found to be of 550 KeV for a $5/2^+$ quasiparticle, mostly due to variations in $\hbar \Omega_{\text{HO}}$. Therefore, we estimate theoretical errors of $\leq 1$ MeV for the Dyson-ADC(3) gaps discussed below. For Gorkov calculations, we expect that errors on binding and separation energies will be larger due to the simpler many-body truncation. However, we note that Ref. [46] has reported a remarkable independence of dominant quasiparticle peaks on the $\lambda_{\text{SRG}}$ cutoff already at second order.

In the following sections, we will report the results obtained for $\hbar \Omega_{\text{HO}} = 24$ MeV and $\lambda_{\text{SRG}} = 2$ fm$^{-1}$.

![FIG. 1. (Color online) Single particle spectral distributions for the addition and removal of a proton to/from closed sub-shell oxygen isotopes. States above the Fermi Surface ($E_F$) are indicated by the shaded areas and yield the spectra of the resulting odd-even fluorine isotopes. The spectra below $E_F$ is for odd-even nitrogen isotopes in the final state (this appears inverted in the plot, with higher excitation energies pointing downward). Fragments with different angular momentum and parity are shown with different colors, as indicated, and the bar lengths provide the calculated spectroscopic factors. These results are obtained from ADC(3) and the full NN+3NF interaction with $\lambda_{\text{SRG}} = 2.0$ fm$^{-1}$.

IV. RESULTS

A. Spectral functions and evolution of single particle spectra

Three dimensional plots of the full spectral function, Eq. (3), are illustrated in Appendix A. Here, we focus on the energy distribution of the spectral strength calculated by its diagonal part over the single particle degrees of freedom,

$$S(\omega) = \sum_{\alpha} S^+_{\alpha \alpha}(\omega) + S^-_{\alpha \alpha}(\omega)$$

$$= \sum_{n} SF_n^+ \delta(\omega - E_n^{A+1} + E_0^{A})$$

$$+ \sum_{n} SF_k^- \delta(\omega - E_0^{A} + E_k^{A-1}),$$

which yields the energy distribution of spectroscopic factors. Each peak represents an eigenstate of a neighbouring odd-even isotope, whose energy is directly observed.
in nucleon addition and removal experiments.

The particle and hole contributions to Eq. (13), calculated with Dyson ADC(3), are displayed in Fig. 1 for protons and in Fig. 2 for neutrons. The nucleon addition part of the spectra are highlighted by the shaded areas. These figures show the general features of the correlated spectral distribution, which conserves strong quasiparticle fragments close the the Fermi surface but becomes heavily fragmented as one moves further away due to coupling to 2p1h and 2h1p (or more complex) excitations. Quasiparticle states with positive energies are above the one-nucleon continuum threshold ($E_n^{A+1} - E_0^{A} = 0$ MeV) and therefore represent states for scattering of a nucleon off the $|\Psi_0^A\rangle$ target. Since we assume a discrete model space in our calculations the associated particle continuum is found discretised in several peaks that become more dense with increasing energy, reflecting the changes in the density of states for $|\Psi_{n}^{A+1}\rangle$. Quasihole fragments at large negative energies correspond to excited $|\Psi_{k}^{A-1}\rangle$ states and also display a continuum portion of the spectrum. However, the spectral strength for nucleon removal is less pronounced. This due to the fewer degrees of freedom available to generate 2h1p configurations, and it can be equivalently explained in terms of the small overlap between the (A-1)-nucleon wave functions in the continuum and the bound $|\Psi_0^A\rangle$ ground state.

Figure 3 shows the details of the evolution of the dominant proton quasiparticle and quasihole peaks, in the $sd$ and $p$ shells, for increasing neutron number. These are corrected for the effects of the COM motion according

![Figure 2](image2.png)

**FIG. 2.** (Color online) Same as Fig. 1 but for the addition and removal of a neutron. Both particle (shaded areas) and hole spectra are for odd-even oxygen isotopes.

![Figure 3](image3.png)

**FIG. 3.** (Color online) Energy evolution of dominant proton quasiparticle fragments around oxygen isotopes of increasing neutron numbers. These level are for protons in the $p$ and $sd$ shells and refer to the ground and excited states of odd-even nitrogen ($1/2^-, 3/2^-$) and fluorine ($5/2^+, 1/2^+, 3/2^+$). Dots (joined by dashed lines) shows ADC(3) results obtained with the induced NN+3NF interaction. Squares (with full lines) refer to the full Hamiltonian, with the leading NNLO-3NF included. These are the dominant peaks also displayed in Fig. 1. In all cases $\lambda_{SRG} = 2.0$ fm$^{-1}$. Experimental values are from Refs. [47–50].

![Figure 4](image4.png)

**FIG. 4.** (Color online) Same as Fig. 3 but for dominant neutron fragments. Dots (joined by dashed lines) shows ADC(3) results obtained with the induced NN+3NF interaction. Squares (with full lines) refer to the full Hamiltonian and the same shown in Fig. 2. Experimental values are from Refs. [47, 48, 51–53].
For neutrons, they are eigenstates of odd-even $A^{+1}$F isotopes indicated in the second row. In the case of neutrons, they are states of odd-even $A^{\pm}O$. Note that for neutrons and $A=22,24$, these two levels are found across the Fermi surface and correspond to eigenstates of different isotopes. Results are reported for both the induced and full Hamiltonians and $\Delta\varepsilon_{dLS} = \varepsilon_{dLS}^{FULL} - \varepsilon_{dLS}^{IND}$ are the changes due to adding the original 3NFs at NNLO. Experimental values are from Refs. [47, 48, 50–53].

\[
\begin{array}{cccccccc}
\varepsilon_{dLS} & \varepsilon_{pLS} & \varepsilon_{pLS}^{IND} & \varepsilon_{pLS}^{FULL} & \varepsilon_{pLS}^{exp.} & \varepsilon_{pLS}^{exp.} & \varepsilon_{pLS}^{exp.} & \varepsilon_{pLS}^{exp.} \\
\varepsilon_{1/2}^{1/2} & \varepsilon_{1/2}^{1/2} & \varepsilon_{1/2}^{1/2} & \varepsilon_{1/2}^{1/2} & \varepsilon_{1/2}^{1/2} & \varepsilon_{1/2}^{1/2} & \varepsilon_{1/2}^{1/2} & \varepsilon_{1/2}^{1/2} \\
15F & 17F & 23F & 25F & 29F & 15O & 17O & 23\pm2O & 25\pm2O & 27O \\
\hline
\varepsilon_{1/2}^{1/2} & 1.32 & 1.67 & 2.72 & 3.02 & 2.92 & 2.70 & 1.77 & 3.06 & 3.30 & 2.03 \\
\varepsilon_{1/2}^{1/2} & 3.17 & 5.33 & 1.17 & 1.84 & 4.98 & 6.33 & 5.65 & 5.05 & 5.06 & 6.28 \\
\varepsilon_{1/2}^{1/2} & 4.48 & 7.00 & 3.88 & 4.86 & 7.90 & 9.02 & 7.42 & 8.12 & 8.36 & 8.32 \\
\varepsilon_{1/2}^{1/2} & 5.00 & 3.83/3.44 & 3.83 & 3.44 & 5.09 & 7.38 & 6.64 & 6.64 & 6.64 \\
\end{array}
\]

TABLE I. Spin-orbit splittings between $1/2^+$ of the odd-even $A^{+1}$F isotopes. For neutrons, these are eigenstates of odd-even $A^{\pm}O$. The evolution of the splitting between corresponding neutron holes, which is close to the empirical value of 6.32 MeV. However, the full Hamiltonian increases it to 8.2 MeV, overestimating the experiment. Exactly the same situation is found for the splitting between corresponding neutron holes, which is also increased by 3.1 MeV due to the original 3NF at NNLO. For comparison, the Argonne v18 interaction that has a strongly repulsive core predicts a separation of $\approx 3.1$ MeV [54] for these two states, at the NN interaction level. The corresponding Urbana IX 3NF increases this by another 2.7 MeV predicting a splitting that is much closer to the experiment [55]. Both three-nucleon Hamiltonians include two-pion terms of the Fujita-Miyazawa type and it is therefore reasonable that they generate similar corrections.

From Tabs. I and II it is clear that the present NN+3N chiral Hamiltonians have a slight tendency to stretch the single particle spectrum, as compared to the experimentally observed dominant peaks. Corrections to this flaws may come at the price of introducing extra short-range repulsion in the NN interactions, for example through higher chiral cutoffs. At lower resolution scales, this implies the possible presence of relevant many-body forces at least at the 4NF level. Ref. [22] pointed out that the

TABLE II. Spin-orbit splitting between $1/2^+$ and $3/2^+$ quasi-particle fragments. For the case of protons these are eigenstates of the odd-even $A^{+1}$N isotopes. For neutrons, these are eigenstates of odd-even $A^{\pm}O$, except for $^{14}$O where the two quasiparticles correspond to different isotopes. Results are reported for both the induced and full Hamiltonians and $\Delta\varepsilon_{pLS} = \varepsilon_{pLS}^{FULL} - \varepsilon_{pLS}^{IND}$ are the contributions due to original 3NFs at NNLO. Experimental values are from Refs. [47, 49].

\[
\begin{array}{cccccccc}
\varepsilon_{pLS} & \varepsilon_{pLS}^{IND} & \varepsilon_{pLS}^{FULL} & \varepsilon_{pLS}^{exp.} & \varepsilon_{pLS}^{exp.} & \varepsilon_{pLS}^{exp.} & \varepsilon_{pLS}^{exp.} & \varepsilon_{pLS}^{exp.} \\
\varepsilon_{1/2}^{3/2} & \varepsilon_{1/2}^{3/2} & \varepsilon_{1/2}^{3/2} & \varepsilon_{1/2}^{3/2} & \varepsilon_{1/2}^{3/2} & \varepsilon_{1/2}^{3/2} & \varepsilon_{1/2}^{3/2} & \varepsilon_{1/2}^{3/2} \\
13N & 15N & 21N & 23N & 27N & 15.13O & 15O \\
\hline
\varepsilon_{1/2}^{3/2} & 0.92 & 0.00 & 2.84 & 4.52 & 3.01 & 0.61 & -2.03 & -2.14 & -3.64 & -1.74 & -2.20 \\
\varepsilon_{1/2}^{3/2} & 2.00 & 3.10 & 2.84 & 4.52 & 3.01 & 0.92 & 0.00 & 2.84 & 4.52 & 3.01 & 0.61 & -2.03 & -2.14 & -3.64 & -1.74 & -2.20 \\
\varepsilon_{1/2}^{3/2} & 0.92 & 0.00 & 2.84 & 4.52 & 3.01 & 0.61 & -2.03 & -2.14 & -3.64 & -1.74 & -2.20 \\
\varepsilon_{1/2}^{3/2} & 2.00 & 3.10 & 2.84 & 4.52 & 3.01 & 0.92 & 0.00 & 2.84 & 4.52 & 3.01 & 0.61 & -2.03 & -2.14 & -3.64 & -1.74 & -2.20 \\
\end{array}
\]

TABLE III. Energy gaps between the dominant $5/2^+$ and $1/2^-$ quasiparticles. These give a measure of the gaps between the $sd$ and $p$ shells for the $^{14}$O isotopes indicated in the second row. For the case of protons, these are particle-hole gaps and coincide with the ground states of the corresponding odd-even $A^{+1}$F and $A^{\pm}O$ isotopes. For neutrons, these are eigenstates of odd-even $A^{+1}$O but are not necessarily situated across the Fermi surface. Results are reported for both the induced and full Hamiltonians and $\Delta E_{gap} = E_{gap}^{FULL} - E_{gap}^{IND}$ are the effects of the original 3NFs at NNLO. Experimental data are from Refs. [47–53].
experimental gaps between the $sd$ and $pf$ shells are over estimated for the Ca and neighbouring isotopic chains. To investigate the behaviour in the present case, we consider the separations between the dominant $5/2^+$ and $1/2^-$ fragments that is representative of the gap between the $p$ and $sd$ shells. These are reported in Tab. III. With the only exception of the proton gap in $^{14}$O, we find that pre-existing 3NFs have the effect of reducing the distance between the two shells, by about 2 MeV, and to bring it closer to the experiment. In spite of this, the gaps are consistently predicted too large by just a few MeV even when the full Hamiltonian is used.

### B. Spectroscopic factors

The quenching of spectroscopic factors (SFs) for the dominant quasiparticle peaks can provide useful insights on the strength of the correlations generated by the Hamiltonian. The principal mechanisms that are responsible for these, are identified with the coupling of nucleons to high-momentum states due to short-range physics and with long-range effects that include collective resonance modes and configuration mixing at small excitation energies [56]. Even for Hamiltonians that present strongly repulsive NN cores, the effects of short-range physics is usually found to be at most a 10% reduction with respect to the independent particle model prediction [1]. Thus, the quenching of SFs is mostly a consequence of low-energy physics. For the present Hamiltonian, the SRG evolution completely removes any quenching due to short-range correlations.

Since spectroscopic factors are directly linked to the cross sections probed by particle addition and removal processes, it has long been debated whether their evolution with proton-neutron asymmetry can explain the observed variations in the strength of direct nucleon knock-out cross sections [57, 58]. The difference between the proton and neutron separation energies is normally taken as a measure of such asymmetry. A case of particular interest is $^{14}$O because of the very large value of this quantity. The present ADC(3) calculations yield substantially the same spectroscopic factors equal to 77.4% (77.2%) for the removal of a proton (neutron) from this isotope to the ground state of $^{13}$N ($^{13}$O). Recent measurements of the $(d,^3$He) and $(d,^3$H) reactions are found to be consistent with our calculations and therefore support a near independence of correlations effects from proton-neutron asymmetry [59].

In order to extend the analysis to cases with larger differences between proton and neutron numbers, we plot in Fig. (5) the SFs for ground state to ground state transitions along the whole chain. In general, we find values evenly spread between 70% and 90% of the independent particle model. The smaller values of SFs are obtained at low separation energies and involve transitions to/from $^{14,22,24}$O. These isotopes present reduced particle-hole neutron gaps and therefore allow for stronger correlations at the Fermi surface. This consideration is also consistent with previous works that clearly showed a close correlations between the particle-hole gap at the Fermi surface and the predicted values of SFs [27]. From this, one may infer that the over diluted spectra reported in Tabs. I, II and III result in more modest quenchings of SFs than otherwise expected.

By looking only at transitions that involve the doubly-closed major shells $^{16}$O and $^{28}$O, one can still identify a correlation between SFs and nucleon separation energies. In particular, proton orbits tend to be more deeply bound as the number of neutrons increases. This is due to the strong components of the proton-neutron forces, which also enhances their correlations. However, the overall dependence on proton-neutron asymmetry is rather mild. We note that the vicinity to the neutron dripline would require to explicitly account for the continuum. Ref. [60] found that this effect is sizeable for $^{24,28}$O and leads to further quenching of the proton SFs.

For the case of the induced Hamiltonian we find a completely similar picture, with SFs of dominant peaks being on average slightly larger than those obtained with the full interaction. Also in this case, stronger quenchings are associated with increased fragmentation of nearby strength and the narrowing of subshell gaps. Thus, we conclude that the general effects of the original 3NFs on the quenching of absolute SFs are indirect and result from the rearrangement of shell orbits and excitation gaps.
C. Results for open shells

The present implementation of the Gorkov-GF approach allows calculations up to the second order in the self-energy [i.e. at the ADC(2) level]. This yields proper predictions of the trend of binding energies [22] but it does not guarantee highly accurate one-nucleon separation energies [40].

We plot the Gorkov predicted binding energies for all even-even isotopes in Fig. 6 and compare them to the Dyson-ADC(3) results where available. For the Dyson case, the induced Hamiltonian systematically under binds the full isotopic chain and predicts $^{20}\text{O}$ to be bound with respect to $^{24}\text{O}$. This is fully corrected by including the original 3NF at leading order, which brings all results to about 3% form the experiment or closer. This is well within the estimated theoretical errors discussed above [19]. The dot-dashed line shows the trend of ground state energies for the full Hamiltonian obtained form Gorkov, which include the $^{18,20,26}\text{O}$ isotopes. This demonstrates that the fraction of binding missed by the second order truncation is rather constant across the whole isotopic chain and, in the present case, of about 2-4 MeV. The result is a constant shift with respect to the complete ADC(3) prediction and the overall trend of binding energy is reproduced very close to the experiment. Note that binding energies for odd-even oxygens can be calculated either as neutron addition or neutron removal from two different nearby isotopes. Fig. 6 shows that this procedure can lead to somewhat different results, which should be taken as an indication of the errors due to the second order many-bod truncation. For the more complete Dyson-ADC(3) method and the full Hamiltonian, these differences are never larger than 200 keV and are not visible in the plot.

Figure 7 shows the analogous information for the binding energies of the nitrogen and fluorine isotopic chains, obtained through removal and addition of one proton. This confirms that all considerations made regarding the effects of leading order 3NFs on the oxygens also apply to their neighbouring chains. In particular, the repulsive effect on the $d_{3/2}$ neutron orbit is key in determining the neutron driplines at $^{23}\text{N}$ and $^{24}\text{O}$. Fluorine isotopes have been observed experimentally up to $^{31}\text{F}$ but with a $^{29}\text{F}$ that is very weakly bound. Fig. 7 clearly demonstrates that this is due to an very subtle cancellation between the repulsion form 3NFs and the attraction generated by one extra proton [19]. Our calculations with the more accurate Dyson-ADC(3) scheme predict $^{29}\text{O}$ to be unbound with respect to $^{24}\text{O}$ by 5.2 MeV. However, this value should be slightly affected by the vicinity to the continuum [17], which was neglected in the present work.

The general qualitative features of the spectral functions discussed in the previous sections are also found in our Gorkov propagators but with an even more dilute single particle spectra. For example, the separation among the $1/2^-$ and $3/2^-$ quasiholes of $^{15}\text{N}$ is found to be 10.2 MeV, compared to the 8.2 MeV calculated in the Dyson-ADC(3) scheme [cfr. Tab. II]. This larger value is a consequence of neglecting the interactions between $2p_{1/2}$ and $2h_{11/2}$ configurations by the second order truncation. Interestingly, this splitting is sensibly reduced in the neighbouring semi-magic isotopes and it is calculated to be 4.9 MeV for $^{17}\text{N}$ and 5.6 MeV for $^{19}\text{N}$. These smaller values are in parts due to the fact the the $p_{3/2}$ hole orbit is fragmented for these nuclei. And we report here only the values for the first $3/2^-$ state.
close to the Fermi surface. In the Gorkov calculations, the changes in these splitting due to aiding the original NNLO-3NF \( \Delta_{pL} \) are 4.2, 0.8 and 2.2 MeV, respectively for \( ^{15}\text{N} \), \( ^{17}\text{N} \) and \( ^{19}\text{N} \).

Figure 8 demonstrates the trend obtained for matter radii along the whole chain and compares them to the observed charge radii for \( ^{16}\text{O} \) and \( ^{18}\text{O} \). The induced 3NFs give sizeable contributions to the calculated radii, which would be sensibly smaller if computed from the evolved NN-only interaction [63]. Conversely, the original 3NFs cause only a small reduction. Eventually, the radii predicted by the complete Dyson-ADC(3) calculations and the full Hamiltonian are smaller than the experimental charge radii by 0.3-0.4 fm. Note that the present calculations do not account for the evolution of operators through the SRG. On the other hand, the small radii are consistent with the dilute spectra discussed above. The radii calculated with the second order Gorkov approach give somewhat smaller results due to the many-body truncation. Nevertheless, they describe the overall increase in matter radii along the whole chain.

V. CONCLUSIONS

We have presented a comprehensive study of the single particle spectral functions and ground state properties of oxygen isotopes based on chiral NN+3N interactions. To this goal, we performed \textit{ab initio} calculations within self-consistent Green’s function (SCGF) theory. The theoretical framework of this approach has been reviewed highlighting the physics information contained in the nuclear spectral function. Calculations were performed for the closed subshell isotopes using the Dyson-ADC(3) many-body truncation scheme which is currently the state of the art technique. For the open shell isotopes \( ^{18}\text{O} \), \( ^{20}\text{O} \) and \( ^{26}\text{O} \) we performed calculations using the recently introduced Grokov formulation of SCGF, which can be currently applied at second order.

The general features of the nuclear spectral functions have been discussed, with particular emphasis on the quasiparticle energies (also referred to as ‘separation’ energies) for the dominant peaks observed in the removal and addition of a nucleon. The 3N interactions at NNLO have the effect of increasing the spin orbit splittings of the \( p \) and \( d \) orbits and lead to over estimating the experiment. At the same time, the 3NFs reduce the gaps between the \( p \) and \( sd \) major shells, improving the agreement with data but not enough to reproduce the empirical values. We observe that all these deficiencies might be corrected by having extra short-range repulsion in the NN section of the Hamiltonian. In general, it is found that the current NN-N3LO interaction with cutoff at 500 MeV, augmented by the local 3NF-NNLO with a 400 MeV local cutoff, tend to dilute the single particle spectrum too much compared to data. The corresponding predictions for matter radii under estimate the experiment.

The conclusion that the present chiral forces overestimate the gaps between major shells was already pointed out in Ref. [22] for isotopic chains around Ca, and suggest that these saturate nuclear matter at slightly higher densities than the empirical point. Here, we find that hints of the same pathologies are seen also for the oxygen isotopes, in spite of the fact that binding energies are nicely predicted at smaller masses.

The calculated absolute quenchings of spectroscopic factors change only mildly with proton-neutron asymmetry. This is valid as long as the occupied states in the single particle spectrum are not near to the continuum. Stronger correlations are instead generated by smaller particle-hole gaps in the oxygen isotopes with closed subshells.

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Appendix A: plots of spectral functions

The diagonal part of the one-nucleon spectral function, Eq. (3), has a straightforward physical interpretation [64, 65]. Its particle part, \( S_{p\alpha}(\omega) \), is the joint proba-
FIG. 9. (Color online) Diagonal part of the complete proton spectral function, Eq. (A1), for closed subshell isotopes $^{14,16,22,24,28}$O. The discretised energy peaks that appear as energy delta functions in Eq. (3) have been smeared with Lorentzians of suitable width. Energies below the Fermi surface, $E_F$, correspond to the hole part of the spectral distribution while those above are for particle addition. The part for $\omega > 0$ MeV (plotted in red) correspond to proton-nucleus scattering states.

FIG. 10. (Color online) Same as Fig. 9 but for neutrons.

The probability of adding a nucleon with quantum numbers $\alpha$ to the A-body ground state, $|\Psi_0^A\rangle$, and then to find the system in a final state with energy $E^{A+1} = E_0^A + \omega$. Likewise, $S^h_{\alpha\alpha}(\omega)$ gives the probability of removing a particle from state $\alpha$ and later finding the nucleus in an eigenfunction of energy $E^{A-1} = E_0^A - \omega$. Once transformed to coordinate or momentum representations, these distributions...
give a rather intuitive picture of the single particle structure of a nucleus. We demonstrate this by calculating the spectral function in coordinate basis as

\[
S(r, \omega) = \sum_{\alpha \beta} \phi_\alpha(r) \left[ S_{\alpha\beta}^p(\omega) + S_{\alpha\beta}^n(\omega) \right] \phi_\beta^*(r), \tag{A1}
\]

where \(\phi_\alpha(r)\) are spin-coupled harmonic oscillator functions. The angular and spin dependences in Eq. (A1) are removed by summing over all oscillator states. This applies for all nuclei under consideration because they have \(J^s = 0^+\) quantum numbers in their ground states.

The spectral functions obtained from Dyson-ADC(3) calculations are displayed in Figs. 9 and 10 for protons and neutrons, respectively. These show the radial distribution of the squared one-nucleon overlap wave functions and neutrons, respectively. These show the radial distributions, which cannot be normalised to the usual asymptotic boundary conditions. Nevertheless, the plots yield the nucleon density \(\rho(r)\) [see Eq. (4)] and further integrations over coordinate space yields the particle number. Note that quasiparticle states for \(\omega > 0\) correspond to the continuum spectrum of the corresponding (A+1)-nucleon system. These are unbound states for the scattering of a nucleon off the \(\{|\Psi^A\rangle\}\) ground state [35, 36]. Thus, they extend to infinity in the limit of a complete single particle model space. In the present work, we only calculate their projection on a truncated harmonic oscillator space, which cannot be normalised to the usual asymptotic boundary conditions. Nevertheless, the plots put in evidence the predicted location for neutron resonances in the \(sd\) and \(pf\) shells. Proper boundary conditions can be obtained by transforming the self-energy, Eq. (11), to coordinate or momentum space. This is normally done in applications of SCGF theory to scattering [66, 67].

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