Approximate relativistic optimized potential method

T. Kreibich and E. K. U. Gross
Institut für Theoretische Physik, Universität Würzburg, Am Hubland, D-97074 Würzburg, Germany

E. Engel
Institut für Theoretische Physik, Universität Frankfurt, Robert Mayer Straße 8-10, D-60054 Frankfurt-am-Main, Germany

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Approximate semianalytical solutions of the integral equation for the relativistic optimized potential are constructed by extending a method recently proposed by Krieger, Li, and Iafrate [Phys. Lett. A 146, 256 (1990)] to the relativistic regime. The quality of the approximation is tested in the longitudinal x-only limit where fully numerical solutions of the relativistic optimized effective potential integral equation are available for spherical atoms. The results obtained turn out to be in excellent agreement with the exact x-only values. The proposed method provides significant improvement over the conventional relativistic local density approximation and generalized gradient approximation schemes. [S1050-2947(97)04912-3]

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I. INTRODUCTION

Since the seminal work of Hohenberg and Kohn [1] and Kohn and Sham [2], density functional theory (DFT) has become a powerful tool for \textit{ab initio} electronic structure calculations of atoms, molecules, and solids [3–5]. The development of more and more refined approximations of the exchange-correlation (xc) energy functional $E_{xc}$ has led to significant improvement over the standard local density approximation (LDA). In particular, the so-called optimized potential method (OPM) [6,7] employing explicitly orbital-dependent functionals rather than the traditional density-dependent functionals has achieved highly accurate results [7–20].

Most of the advances in DFT have been made in the context of nonrelativistic physics. For high-Z atoms, however, relativistic contributions have to be considered. For example, the ground-state energy of mercury (Hg) decreases from its nonrelativistic value $\approx 18,408$ a.u. to $\approx 19,649$ a.u. if relativistic effects are taken into account. Furthermore, even for systems with moderate Z, relativistic contributions to $E_{xc}$ are larger than the differences between the currently best xc energy functionals [21]. Until now, the calculation of such relativistic contributions has mostly been based on the relativistic local density approximation (RLDA). To go beyond the RLDA, an x-only version of the OPM was formulated for relativistic systems [22,21]. As in the nonrelativistic case, the solution of the resulting equations is a rather demanding task and has been achieved so far only for systems of high symmetry, e.g., spherical atoms [21,23–25].

The purpose of the present paper is to develop a simplified version of the relativistic OPM (ROPM) scheme leading to a generalization of the approximation of Krieger, Li, and Iafrate (KLI) [8–10,26–36] to the realm of relativistic systems. This will be done for systems subject to arbitrary static external four-potentials. The paper is organized as follows. In Sec. II we give a brief review of the foundations of relativistic DFT (RDFT). After that, in Sec. III, we develop the ROPM generalizing a nonrelativistic derivation of Görling and Levy [37] to the relativistic domain. The relativistic KLI (RKLI) scheme is developed in Sec. IV before some limiting cases are discussed in Sec. V. In Sec. VI, numerical results of the ROPM and RKLI methods are presented and compared to other RDFT methods.

II. THEORETICAL BACKGROUND

On its most general level, RDFT is based on quantum electrodynamics (QED) and thus contains not only relativistic but also radiative effects. For a detailed derivation, also including questions of renormalization, the reader is referred to recent reviews [23,38]. The central statement of RDFT — the relativistic version of the Hohenberg-Kohn (HK) theorem [39] — can be stated in the following way: The renormalized ground-state four-current $j^0(r)$ of an interacting system of Dirac particles uniquely determines, up to gauge transformations, the external four-potential $A^\mu_{ext}(r)$ as well as the ground-state wave function $\Psi[j^\mu]$. As a consequence, any observable of the relativistic many-body system under consideration is a functional of its ground-state four-current. As in the nonrelativistic case, the exact ground-state four-current including all quantum electrodynamical effects can in principle be obtained from an auxiliary noninteracting system—the relativistic Kohn-Sham (RKS) system [21,38,40,41]:

$$j^0(r) = \sum_{-\epsilon < \epsilon_k < \epsilon_F} \bar{\psi}_k(r) j^\nu(\gamma^\nu \psi_k(r)) + j^\nu_{\nu}(r),$$

where $j^\nu_{\nu}(r)$ denotes the vacuum contribution to the four-current. The four-component spinors $\psi_k(r)$ are solutions of an effective single-particle Dirac equation (in atomic units $\hbar = e = m = 1$)

$$\gamma_\nu[-ic \gamma^\nu \nabla + c^2 + \gamma_\mu A^\mu_{ext}(r)]\psi_k(r) - \epsilon_k \psi_k(r),$$

with $A^\mu_{ext}(r)$ being the effective four-potential, which can be decomposed according to

$$A^\mu_{ext}(r) = A^\mu_{ext}(r) + \int d^4r' \frac{j^\mu(r')}{|r-r'|} + A^\mu_{ext}(r'),$$

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Here, \( A_{\text{xc}}^\mu(r) \) is a static but otherwise arbitrary external four-potential, the second term on the right-hand side represents a Hartree-like potential, and the last term, defined by

\[
A_{\text{xc}}^\mu[j^\nu](r) = \frac{\delta E_{\text{xc}}[j^\nu]}{\delta j^\mu(r)}
\]

(4)
denotes the xc four-potential containing by construction all nontrivial many-body effects.

Equations (1)–(4) represent the relativistic KS scheme that has to be solved self-consistently. However, the calculation of the vacuum contribution \( j^\nu(r) \) to the four-current requires the knowledge of an infinite number of (positive and negative energy) states, so that one would have to deal with an infinite system of coupled equations. Since such a procedure is highly impractical, we will, in the following, ignore all vacuum contributions to the various energy components and to the four-current, which is then given just by the first term on the right-hand side of Eq. (1). This means that we restrict ourselves to the calculation of relativistic effects and neglect radiative corrections. Since we are aiming at electronic structure calculations for atoms, molecules, and solids, we expect the neglected terms to be small. If one were after all interested in the radiative contributions, an \textit{a posteriori} perturbative treatment should be sufficient and represents in fact the standard approach.

III. RELATIVISTIC OPTIMIZED POTENTIAL METHOD

In order to derive a relativistic generalization of the OPM integral equation, we start out from the total-energy functional of a system of interacting Dirac particles (neglecting vacuum contributions) subject to a static external four-potential \( A_{\text{ext}}^\mu(r) \):

\[
E_{\text{tot}}^{\text{ROPM}}[j^\nu] = \sum_{-\epsilon^2 < k^\mu < \epsilon_F} \int d^3r \varphi^*_\mu(r)(-i\epsilon \gamma \cdot \nabla + c^2)\varphi_\mu(r) + \int d^3r j^\mu(r)A_{\text{ext}}^\mu(r)
+ \frac{1}{2} \int d^3r d^3r' \frac{j^\mu(r)j^\nu(r')}}{|r-r'|} + E_{\text{xc}}^{\text{ROPM}}[\varphi_\mu].
\]

(5)

In contrast to the ordinary RDFT approach, the xc energy functional is given here as an explicit functional of the RKS four-spinors \( \varphi_\mu \). Still, \( E_{\text{xc}}^{\text{ROPM}}[\varphi_\mu] \) represents a functional of the density: Via the HK theorem applied to noninteracting systems, \( j^\nu(r) \) uniquely determines the effective potential \( A_{\text{xc}}^\mu[j^\nu] \). With this very potential, the Dirac equation (2) is solved to obtain the set of single-particle orbitals \( \varphi_\mu[j^\nu] \), which are then used to calculate the quantity \( E_{\text{xc}}^{\text{ROPM}}[\varphi_\mu[j^\nu]] \). Therefore every functional, depending \textit{explicitly} on RKS spinors, is an \textit{implicit} functional of the density, provided the orbitals come from a local potential. This allows us to use the exact expression for the longitudinal exchange energy, i.e., the relativistic Fock term

\[
E_{\text{xc}}^{\text{exact}}[j^\nu] = -\frac{1}{2} \sum_{-\epsilon^2 < k^\mu < \epsilon_F} \int d^3r \int d^3r' \times \frac{\varphi^*_\mu(r)\varphi_\mu(r)\varphi^*_\nu(r')\varphi_\nu(r')}{|r-r'|}. 
\]

(6)

One major advantage of such an exact treatment of the exchange energy lies in the fact that the spurious self-interactions contained in the Hartree energy are fully canceled. The price to be paid for the orbital dependence of \( E_{\text{xc}} \) is that the calculation of the xc four-potential \( A_{\text{xc}}^\mu(r) \), defined by Eq. (4), is somewhat more complicated. It has to be determined by an integral equation, as will be shown in the following.

Starting from the definition of \( A_{\text{xc}}^\mu(r) \), Eq. (4), we can calculate the xc four-potential corresponding to an orbital-dependent xc energy functional by applying the chain rule for functional derivatives:

\[
A_{\text{xc}}^{\text{ROPM}}(r) = -\sum_{-\epsilon^2 < k^\mu < \epsilon_F} \int d^3r' \int d^3r'' \times \left( \frac{\delta E_{\text{xc}}^{\text{ROPM}}[\varphi_\mu]}{\delta \varphi_\mu(r')} \frac{\delta \varphi_\mu(r)}{\delta A_{\text{sp}}(r')} + \text{c.c.} \right) \frac{\delta A_{\text{sp}}(r')}{\delta j^\mu(r)}. 
\]

(7)

The last term on the right-hand side of Eq. (7) is readily identified with the inverse of the static response function of a system of noninteracting Dirac particles

\[
\chi_{\text{sp}}^{-1}(r,r') = \frac{\delta j^\mu(r)}{\delta A_{\text{sp}}(r')},
\]

(8)

so that Eq. (9) can be rewritten as

\[
A_{\text{xc}}^{\text{ROPM}}(r) = -\sum_{-\epsilon^2 < k^\mu < \epsilon_F} \int d^3r' \int d^3r'' \times \left( \frac{\delta E_{\text{xc}}^{\text{ROPM}}[\varphi_\mu]}{\delta \varphi_\mu(r')} \frac{\delta \varphi_\mu(r)}{\delta A_{\text{sp}}(r')} + \text{c.c.} \right) \chi_{\text{sp}}^{-1}(r',r).
\]

(9)

Acting with the response operator (8) on Eq. (9) and using the identity

\[
\int d^3r \chi_{\text{sp}}^{-1}(r,r')\chi_{\text{sp}}^\mu(r,r') - \delta_{\nu\sigma} \delta(r-r')
\]

we obtain (after rearranging the indices)

\[
\int d^3r' A_{\text{xc}}^{\text{ROPM}}(r')\chi_{\text{sp}}^\mu(r',r) = \sum_{-\epsilon^2 < k^\mu < \epsilon_F} \int d^3r \frac{\delta E_{\text{xc}}^{\text{ROPM}}[\varphi_\mu]}{\delta \varphi_\mu(r')} \frac{\delta \varphi_\mu(r)}{\delta A_{\text{sp}}(r')} + \text{c.c.}
\]

(11)

To further evaluate this equation, we note that the first functional derivative on the right-hand side of Eq. (11) is readily
computed once an expression for $E_{xc}^{\text{ROPM}}$ in terms of single-particle spinors is given. The remaining functional derivative is calculated by using standard first-order perturbation theory, yielding

$$\frac{\delta \varphi_{\mu}(r')}{\delta A_{sk}(r)} = \sum_{l' k} \frac{\varphi_{\mu}(r')}{\epsilon_{k} - \epsilon_{l'}} \bar{\varphi}(r) \gamma^\mu \varphi_{\mu}(r). \tag{12}$$

This equation also enables us to give an explicit expression for the response function

$$\chi_{sk}^{\mu\nu}(r, r') = -\frac{\delta}{\delta A_{sk}(r')} \left( \sum_{l' k} \bar{\varphi}(r') \gamma^\mu \varphi(r') \gamma^\nu \varphi_{\mu}(r) \right) \tag{13}$$

in terms of the RKS spinors:

$$\chi_{sk}^{\mu\nu}(r, r') = \sum_{l' k} \sum_{-\epsilon_{k} < \epsilon_{l'} < \epsilon_{F}} \bar{\varphi}(r') \gamma^\mu \varphi(r') \gamma^\nu \varphi_{\mu}(r) \frac{\epsilon_{l'}}{\epsilon_{k} - \epsilon_{l'}} + \text{c.c.} \tag{14}$$

Finally, putting Eqs. (11), (12), and (14) together leads to the ROPM integral equations for the local xc four-potential $A_{xc}^{\text{ROPM}}(r)$:

$$\sum_{l' k} \int d^3 r' \left( \bar{\varphi}(r') \gamma^\mu \varphi(r') \gamma^\nu A_{xc}^{\text{ROPM}}(r') \right)$$

$$\quad - \frac{\delta E_{xc}^{\text{ROPM}}}{\delta \varphi_{\mu}(r')} \left( G_{sk}(r', r) \gamma^\mu \varphi_{\mu}(r) \right) + \text{c.c.} = 0, \tag{15}$$

where $G_{sk}(r', r)$ is defined as

$$G_{sk}(r', r) = \sum_{l' k} \frac{\varphi_{\mu}(r') \varphi_{\mu}(r)}{\epsilon_{l'} - \epsilon_{k}}. \tag{16}$$

Now the ROPM scheme is complete: For a given approximation of the xc energy, the ROPM integral equations have to be solved for $A_{xc}^{\text{ROPM}}(r)$ simultaneously with the RKS equation (2) until self-consistency is achieved. Note that Eq. (15) determines the xc four-potential $A_{xc}^{\text{ROPM}}(r)$ only up to an arbitrary constant, which can be specified by requiring $A_{xc}^{\text{ROPM}}(r)$ to vanish asymptotically (for finite systems).

To conclude this section, we note that exchange and correlation contributions can be treated separately within the ROPM scheme. This is most easily seen by starting out with only the exchange four-potential, defined as $A_{xc}^e(r) = \delta E_{xc}/\delta j_{\mu}(r)$, instead of Eq. (4) and repeating the steps which lead to Eq. (15) and likewise for the correlation potential.

**IV. TRANSFORMATION OF THE ROPM INTEGRAL EQUATIONS AND THE RELATIVISTIC KLI APPROXIMATION**

In order to use the ROPM equations derived in the preceding section, we have to solve Eq. (15) for the xc four-potential. Unfortunately, there is no known analytic solution for $A_{xc}^{\text{ROPM}}(r)$ depending explicitly on the set of single-particle spinors $\{\varphi_{\mu}\}$. We therefore have to deal with Eq. (15) numerically, which is a rather demanding task. Thus a simplified scheme for the calculation of $A_{xc}^{\text{ROPM}}(r)$ appears highly desirable.

To this end we shall perform a transformation of the ROPM integral equations similar to the one recently introduced by KLI in the nonrelativistic domain [29,31]. This will lead to an alternative but still exact form of the ROPM integral equation which naturally lends itself as a starting point for systematic approximations. We start out by defining

$$\psi^+(r) = \int d^3 r' \left( \bar{\varphi}(r') \gamma^\mu A_{xc}^{\text{ROPM}}(r') \right)$$

$$\quad - \frac{\delta E_{xc}^{\text{ROPM}}}{\delta \varphi_{\mu}(r')} \left( G_{sk}(r', r) \gamma^\mu \varphi_{\mu}(r) \right) + \text{c.c.} = 0, \tag{17}$$

such that the ROPM integral equations can be rewritten as

$$\sum_{-\epsilon_{k} < \epsilon_{l'} < \epsilon_{F}} \bar{\varphi}(r) \gamma^\mu \varphi_{\mu}(r) + \text{c.c.} = 0, \tag{18}$$

where the adjoint spinor $\bar{\psi}(r)$ is defined in the usual way, i.e.,

$$\bar{\psi}_{\mu}(r) := \varphi^+_{\mu}(r) \gamma^\mu. \tag{19}$$

Since the RKS spinors $\{\varphi_{\mu}\}$ span an orthonormal set, one readily proves the orthogonality relation

$$\int d^3 r \bar{\psi}^+_{\mu}(r) \varphi_{\nu}(r) = 0. \tag{20}$$

We now use the fact that $G_{sk}(r', r)$ is the Green function of the RKS equation projected onto the subspace orthogonal to $\varphi_{\mu}(r)$, i.e., it satisfies the equation

$$G_{sk}(r', r)[\hat{h}^+_{sk}(r) - \epsilon_{L}] = -[\delta(r' - r) - \varphi_{\mu}(r') \varphi^+_\mu(r)], \tag{21}$$

The operator $\hat{h}^+_{sk}(r)$ denotes the Hermitian conjugate of the RKS Hamiltonian,

$$\hat{h}^+_{sk}(r) := - \gamma^\mu \left[ ic \gamma \cdot \hat{v} + c^2 + \gamma^\nu A_{sk}(r) \right]. \tag{22}$$

acting from the right on the unprimed variable of $G_{sk}(r', r)$ (the arrow on top of the gradient indicates the direction in which the derivative has to be taken). Using Eq. (21), we can act with the operator $[\hat{h}^+_{sk}(r) - \epsilon_{L}]$ from the right on Eq. (17), leading to
\[ \psi_k^+(\mathbf{r})[\hat{h}_D^+(\mathbf{r}) - \epsilon_k] = -\left( \frac{\varphi_k^+(\mathbf{r})}{\gamma A_{\text{xc}}^+(\mathbf{r}) - \frac{\delta E_{\text{xc}}}{\delta \varphi_k^+(\mathbf{r})}} \right) + \varphi_k^+(\mathbf{r})(\hat{A}_{\text{xc}}^\text{R(PM)} - \tilde{u}_{\text{xc}}), \] (23)

where we have introduced \( \hat{A}_{\text{xc}}^\text{R(PM)} \) as a shorthand notation for the average value of \( \gamma' A_{\text{xc}}^\text{R(PM)}(\mathbf{r}) \) with respect to the \( k \)th orbital, i.e.,

\[ \hat{A}_{\text{xc}}^\text{R(PM)} : = \int d^3r \varphi_k^+(\mathbf{r})\gamma' A_{\text{xc}}^\text{R(PM)}(\mathbf{r})\varphi_k(\mathbf{r}) \] (24)

and

\[ \tilde{u}_{\text{xc}} : = \int d^3r \frac{\delta E_{\text{xc}}}{\delta \varphi_k^+(\mathbf{r})} \varphi_k^+(\mathbf{r}) \]. (25)

The differential equation (23) has the structure of a RKS equation with an additional inhomogeneity term. Together with the boundary condition

\[ \psi_k^+(\mathbf{r}) \to 0, \] (26)

Eq. (23) uniquely determines \( \psi_k^+(\mathbf{r}) \). To prove this statement, we assume that there are two independent solutions of Eq. (23), namely, \( \psi_k^+(\mathbf{r}) \) and \( \psi_k^\dagger(\mathbf{r}) \). Then the difference of these two solutions, \( \Psi_k^+(\mathbf{r}) : = \psi_k^+(\mathbf{r}) - \psi_k^\dagger(\mathbf{r}) \), satisfies the homogeneous RKS equation

\[ \Psi_k^+(\mathbf{r})[\hat{h}_D^+(\mathbf{r}) - \epsilon_k] = 0, \] (27)

which has a unique solution

\[ \Psi_k^+(\mathbf{r}) = \varphi_k^+(\mathbf{r}), \] (28)

if the above boundary condition is fulfilled. However, this solution leads to a contradiction to the orthogonality relation (20) so that \( \Psi_k^+(\mathbf{r}) \) can only be the trivial solution of Eq. (27),

\[ \Psi_k^+(\mathbf{r}) = 0, \] (29)

which completes the proof.

As an interesting aside, we briefly consider the physical meaning of the quantity \( \psi_k^+(\mathbf{r}) \). Defining

\[ u_{\text{xc}}^{\alpha\beta}(\mathbf{r}) : = \frac{\delta \alpha}{\delta \varphi_{\alpha}^\dagger(\mathbf{r})} \frac{\delta E_{\text{xc}}}{\delta \varphi_{\beta}(\mathbf{r})}, \] (30)

with \( \alpha \) and \( \beta \) denoting spinor indices running from 1 to 4, we can rewrite Eq. (17) as

\[ \psi_k^+(\mathbf{r}) = \sum_{\epsilon_k} [\varphi_{\epsilon_k}^+(\mathbf{r})]^{\alpha\beta} u_{\text{xc}}^{\alpha\beta}(\mathbf{r}) \left( \gamma^0 \gamma^\nu A_{\text{xc}}^\text{R(PM)}(\mathbf{r}') \right) - \tilde{u}_{\text{xc}}(\mathbf{r}) \varphi_1^\dagger(\mathbf{r}'), \] (31)

where summation over \( \alpha \) and \( \beta \) is implicitly understood.

From this equation, it is obvious that \( \psi_k^+(\mathbf{r}) \) is the usual first-order shift in the wave function caused by a perturbing potential \( \delta \text{xc} = (\gamma^0 \gamma^\nu A_{\text{xc}}^\text{R(PM)} - \tilde{u}_{\text{xc}}) \). This fact also motivates the boundary condition assumed above. In \( x \)-only theory, \( \tilde{u}_{\text{xc}}(\mathbf{r}) \) is the local, orbital-dependent RHF exchange potential so that \( \psi_k^+(\mathbf{r}) \) is the first-order shift of the RKS wave function towards the RHF wave function. However, one has to realize that the first-order shift of the orbital-dependent quantity \( \tilde{u}_{\text{xc}}[\{\varphi_k\}] \) has been neglected.

Now we use Eq. (23) to further transform the ROPM equations (18). As a first step, we solve Eq. (23) for \( A_{\text{xc}}^0(\mathbf{r})\Psi_k^+(\mathbf{r}) \), leading to

\[ A_{\text{xc}}^0(\mathbf{r})\psi_k^+(\mathbf{r}) = -\left( \frac{\varphi_k^+(\mathbf{r})}{\gamma A_{\text{xc}}^\text{R(PM)}(\mathbf{r}) - \frac{\delta E_{\text{xc}}}{\delta \varphi_k^+(\mathbf{r})}} \right) + (\hat{A}_{\text{xc}} - \tilde{u}_{\text{xc}})\varphi_k^+(\mathbf{r}) - \tilde{u}_{\text{xc}}(\mathbf{r})[ic \gamma \cdot \nabla + c^2 - \gamma \cdot A_{\text{xc}}(\mathbf{r}) - \gamma^0 \epsilon_k]. \] (32)

We then multiply the ROPM equations (18) by the zeroth component of the effective RKS four-potential, \( A_{\text{xc}}^0(\mathbf{r}) \), yielding

\[ \sum_{\epsilon_k < \epsilon_F \leq \epsilon_F} A_{\text{xc}}^0(\mathbf{r})\tilde{u}_{\text{xc}}(\mathbf{r}) \varphi_k^+(\mathbf{r}) \gamma^\mu \psi_k(\mathbf{r}) + \text{c.c.} = 0, \] (33)

and employ Eq. (32) to obtain

\[ \sum_{\epsilon_k < \epsilon_F \leq \epsilon_F} \left( \varphi_k^+(\mathbf{r}) \gamma^\nu A_{\text{xc}}^\text{R(PM)}(\mathbf{r}) - \frac{\delta E_{\text{xc}}}{\delta \varphi_k^+(\mathbf{r})} \right) + \tilde{u}_{\text{xc}}(\mathbf{r}) \varphi_k^+(\mathbf{r}) \right] + \tilde{u}_{\text{xc}}(\mathbf{r})[ic \gamma \cdot \nabla + c^2 - \gamma \cdot A_{\text{xc}}(\mathbf{r}) - \gamma^0 \epsilon_k] \varphi_k^+(\mathbf{r}) + \text{c.c.} = 0. \] (34)

Introducing the 4X4 matrix

\[ J^{\mu\nu}(\mathbf{r}) : = \frac{1}{2} \sum_{\epsilon_k < \epsilon_F \leq \epsilon_F} \varphi_k^+(\mathbf{r}) \gamma^\nu \gamma^\mu \varphi_k(\mathbf{r}) + \text{c.c.} \] (35)

and defining

\[ a_{\text{xc}}^{\alpha\beta}(\mathbf{r}) : = -\frac{\delta E_{\text{xc}}}{\delta \varphi_k^+(\mathbf{r})} \gamma^0 \gamma^\mu \varphi_k(\mathbf{r}) \] (36)

we rewrite Eq. (34) as

\[ J^{\mu\nu}(\mathbf{r}) A_{\text{xc}}^\text{R(PM)}(\mathbf{r}) = \sum_{\epsilon_k < \epsilon_F \leq \epsilon_F} \left( a_{\text{xc}}^{\alpha\beta}(\mathbf{r}) + j_k^{\alpha\beta}(\mathbf{r}) \hat{A}_{\text{xc}}^\text{R(PM)} - \tilde{u}_{\text{xc}}(\mathbf{r}) \right) \varphi_1^\dagger(\mathbf{r}'), \] (37)

with \( j_k^{\alpha\beta}(\mathbf{r}) \) being the four-current with respect to the \( k \)th orbital.
\[ j_\mu'(\mathbf{r}) = -\vec{\varphi}_\mu'(\mathbf{r}) \gamma^{\mu} \varphi_\mu'(\mathbf{r}). \]  

(38)

In order to solve Eq. (37) for \( A_{x_{\mu}}^{\text{ROP}}(\mathbf{r}) \), we first have to investigate whether the \( 4\times4 \) matrix \( \mathcal{J}(\mathbf{r}) \), defined by Eq. (35), is nonsingular, i.e., whether the inverse \( \mathcal{J}^{-1}(\mathbf{r}) \) exists. We therefore calculate the determinant of \( \mathcal{J}(\mathbf{r}) \), yielding

\[ \det[\mathcal{J}(\mathbf{r})] = [j_0^\mu(\mathbf{r})]^2 j_\mu(\mathbf{r}) = n^4(\mathbf{r}) \left( 1 - \frac{j_0^\mu(\mathbf{r})}{c^2 n^4(\mathbf{r})} \right). \]

(39)

where the last equality follows from the decomposition of the four-current into the scalar density and the vector components according to

\[ j_\mu(\mathbf{r}) = \left( n(\mathbf{r}), \frac{1}{c} \mathbf{j}(\mathbf{r}) \right). \]

(40)

Since the current \( \mathbf{j}(\mathbf{r}) \) divided by the density \( n(\mathbf{r}) \) is the velocity field of the system, it follows from \( v(\mathbf{r}) < c \) that

\[ \det[\mathcal{J}(\mathbf{r})] \neq 0 \]

(41)

and therefore that the matrix \( \mathcal{J}(\mathbf{r}) \) is nonsingular.

Solving then Eq. (37) for \( A_{x_{\mu}}^{\text{ROP}}(\mathbf{r}) \) yields

\[
A_{x_{\mu}}^{\text{ROP}}(\mathbf{r}) - \frac{1}{2} \mathcal{J}^{-1}_{\mu\nu}(\mathbf{r}) \sum_{-c^2 < s_{\nu} < c^2} [a_{x_{\nu}}^{\mu}(\mathbf{r})] \\
+ j_\nu' \left( A_{x_{\nu}}^{\text{ROP}} - \vec{u}_{x_{\nu}} \right) ] + \text{c.c.}
\]

(42)

We emphasize that Eq. (42) is an exact transformation of the ROPM equation (15). In particular, Eq. (42) is still an integral equation. However, its advantage lies in the fact that it naturally lends itself as a starting point for deriving systematic approximations of \( A_{x_{\nu}}^{\text{ROP}}(\mathbf{r}) \). We only need approximate \( \psi_{\mu}^{\nu}(\mathbf{r}) \) in Eq. (36) by a suitable function of the set of RKS orbitals \( \{\varphi_i\} \). The simplest possible approximation is obtained by completely neglecting the terms involving \( \psi_{\mu}^{\nu}(\mathbf{r}) \) in Eq. (36). Although this approximation may seem to be rather crude, it was shown to produce highly accurate results in the nonrelativistic case [29,31].

The xc potential \( A_{x_{\mu}}^{\text{KLI}}(\mathbf{r}) \) is then approximately determined by the following equation:

\[
A_{x_{\mu}}^{\text{KLI}}(\mathbf{r}) = \frac{1}{2} \mathcal{J}^{-1}_{\mu\nu}(\mathbf{r}) \sum_{-c^2 < s_{\nu} < c^2} [ \delta_{\mu\nu} r_{x_{\nu}} ] (\mathbf{r}) \]

\[ + j_\nu' \left( A_{x_{\nu}}^{\text{KLI}} - \vec{u}_{x_{\nu}} \right) \]  

(43)

This equation establishes the generalization of the nonrelativistic KLI approximation to the realm of relativistic systems.

In contrast to the ROPM equation (15), the relativistic KLI (RKLI) equation, although still being an integral equation, can be solved explicitly in terms of the RKS spinors \( \{\varphi_i\} \): Multiplication of Eq. (43) with \( j_\nu' \) \( (\mathbf{r}) \), summing over all \( \mu \), and integrating over space yields

\[ \rho_{x_{\nu}} = \rho_{x_{\nu}}^{\text{KLI}} + \sum_{-c^2 < s_{\nu} < c^2} M_{\mu\nu} \left[ A_{x_{\nu}}^{\text{KLI}} - \frac{1}{2} (\vec{u}_{x_{\nu}} - \vec{u}_{x_{\nu}}^{*}) \right]. \]

(44)

where we have defined

\[ A_{x_{\nu}}^{\text{KLI}} := \int d^3 r j_\mu'(\mathbf{r}) \mathcal{H}_{\mu\nu}(\mathbf{r}) j_\nu'(\mathbf{r}). \]

(46)

The unknown coefficients \( [A_{x_{\nu}}^{\text{KLI}} - \frac{1}{2} (\vec{u}_{x_{\nu}} - \vec{u}_{x_{\nu}}^{*})] \) are then determined by the linear equation

\[ \sum_{-c^2 < s_{\nu} < c^2} (\delta_{\mu\nu} - M_{\mu\nu}) \left[ A_{x_{\nu}}^{\text{KLI}} - \frac{1}{2} (\vec{u}_{x_{\nu}} - \vec{u}_{x_{\nu}}^{*}) \right] = \left[ \rho_{x_{\nu}} - \frac{1}{2} (\vec{u}_{x_{\nu}} - \vec{u}_{x_{\nu}}^{*}) \right]. \]

(47)

Solving this equation and substituting the result into Eq. (43) finally leads to an expression for the xc four-potential \( A_{x_{\nu}}^{\text{KLI}}(\mathbf{r}) \) that depends explicitly on the set of single-particle spinors \( \{\varphi_i\} \). We thus have obtained a method of calculating the xc four-potential \( A_{x_{\mu}}^{\text{ROP}}(\mathbf{r}) \) in an approximate way, which is numerically much less involved compared to the full solution of the ROPM integral equations.

V. ELECTROSTATIC LIMIT

The ROPM and RKLI methods, developed in the preceding sections, can be applied to systems subject to arbitrary static external four-potentials. In particular, the methods allow us to deal with external magnetic fields of arbitrary strength. Yet, in electronic structure calculations of atoms, molecules, and solids, we most commonly encounter situations, where no magnetic fields are present (in a suitable Lorentz frame, typically the rest frame of the nuclei).

Thus in this section we consider four-potentials whose spatial components vanish, i.e., \( A_{x_{\mu}}^{\text{OPM}}(\mathbf{r}) = 0 \). (This also includes a partial fixing of the gauge.) In this situation, a simplified Hohenberg-Kohn-Sham scheme can be developed, stating that the zeroth component \( n(\mathbf{r}) = j_0^\mu(\mathbf{r}) \) of the ground-state current density alone determines the external potential \( V_{\text{ext}}(n) \) and the ground-state wave function \( \Psi[n] \) uniquely (for a discussion on this so-called “electrostatic case” cf. Refs. [21,23]). Consequently, only a scalar effective potential \( V_{\text{ext}}(\mathbf{r}) \) is present in the RKS equation (2).

When orbital-dependent functionals are used for the xc energy in this context, the corresponding scalar xc potential \( V_{\text{xc}}(\mathbf{r}) \) can be calculated by repeating the steps of Sec. III. One then finds the ROPM integral equation for the “electrostatic case”:

\begin{align*}
A_{x_{\nu}}^{\text{OPM}}(\mathbf{r}) & = \frac{1}{2} \mathcal{J}^{-1}_{\mu\nu}(\mathbf{r}) \sum_{-c^2 < s_{\nu} < c^2} [ \delta_{\mu\nu} r_{x_{\nu}} ] (\mathbf{r}) \\
& + j_\nu' \left( A_{x_{\nu}}^{\text{OPM}} - \vec{u}_{x_{\nu}} \right) \]  

(45)

\[
\sum_{-e^2 < \varepsilon < e_F} \int d^3r \left( \psi_\uparrow^\dagger(r') V_{xc}^\text{ROPM}(r') - \frac{\delta E_{\text{xc}}^\text{ROPM}}{\delta \varphi_\downarrow(r')} G_{SS}(r',r) \varphi_\downarrow(r) \right) + \text{c.c.} = 0. \tag{48}
\]

We mention that the same result is obtained if one demands that \( V_\downarrow(r) \) be the variationally best local effective potential yielding single-particle spinors minimizing the total-energy functional (5), i.e.,

\[
\left. \frac{\delta E_{\text{tot}}^\text{ROPM}}{\delta V_\downarrow(r)} \right|_{V_\downarrow=V_\downarrow^\text{ROPM}} = 0. \tag{49}
\]

In fact, using this approach, Shadwick, Talman, and Norman [22] derived the \( x \)-only limit of the ROPM integral equation (48).

Compared to the four ROPM integral equations (15), which determine the \( xc \) four-potential \( J_{\text{xc}}^\text{ROPM} \), Eq. (48) is considerably simpler. Still, its numerical solution is a rather demanding task which has been achieved so far only for systems of high symmetry, i.e., spherical atoms [21–24]. Again, an approximate ROPM scheme can be derived: Following the arguments of Sec. IV, the ROPM integral equation (48) can be exactly rewritten as

\[
V_{xc}^\text{ROPM}(r) = \frac{1}{2n(r) - e^2 < e < e_F} \sum_{\varepsilon < e} n_\downarrow(r) \left[ v_{xc}^* + \left( V_{xc}^\text{ROPM} - \bar{u}_{xc} \right) \right] + \text{c.c.,} \tag{50}
\]

where

\[
v_{xc}^\text{ROPM}(r) = \frac{1}{n_\downarrow(r)} \left( \frac{\delta E_{\text{xc}}^\text{ROPM}}{\delta \varphi_\downarrow(r)} \varphi_\downarrow(r) - i e \nabla \left[ \psi_\downarrow(r) \gamma \varphi_\downarrow(r) \right] \right) \tag{51}
\]

and

\[
\psi_\uparrow^\dagger(r) = \int d^3r' \left( \psi_\uparrow^\dagger(r') V_{xc}^\text{ROPM}(r') - \frac{\delta E_{\text{xc}}^\text{ROPM}}{\delta \varphi_\uparrow(r')} G_{SS}(r',r) \right) \tag{52}
\]

similar to Eq. (17). Equation (50) represents the \( \text{‘electrostatic case’} \) analog of Eq. (42) and can also be approximated in the same way, namely, by neglecting all terms involving \( \psi_\uparrow^\dagger(r) \) in Eq. (51).

In the context of the \( \text{‘electrostatic case’} \) considered here, some more insight into the nature of this approximation can be gained: It can be interpreted as a \( \text{‘mean-field’} \)-type approximation in the sense that the average of the neglected terms with respect to the ground-state density vanishes. To demonstrate this, we note that the neglected terms averaged over \( n(r) \) are given by

\[
\sum_{-e^2 < \varepsilon < e_F} i e \int d^3r \nabla \left[ \psi_\downarrow(r) \gamma \varphi_\downarrow(r) \right] + \text{c.c.} \tag{53}
\]

Applying the divergence theorem, this integral can be transformed to a surface integral which vanishes for finite systems if the surface is taken to infinity. Hence, neglecting the terms involving \( \psi_\uparrow^\dagger(r) \) means replacing them by their average value, which is zero.

The \( xc \) potential \( V_{xc}^\text{ROPM}(r) \) can therefore approximately be determined by the following equation, leading to the RKLI equation for the \( \text{‘electrostatic case’} \):

\[
V_{xc}^\text{RKLI}(r) = \frac{1}{2n(r) - e^2 < e < e_F} \sum_{\varepsilon < e} \left( \frac{\delta E_{\text{xc}}^\text{ROPM}}{\delta \varphi_\downarrow(r)} \right) + n_\downarrow(r) \left( \varphi_{\text{RKLI}}^\text{ROPM} - \bar{u}_{xc} \right) + \text{c.c.} \tag{54}
\]

From this form it is obvious that the RKLI potential closely resembles the relativistic Dirac-Slater potential as well as the nonrelativistic KLI potential. Whether the accuracy of the corresponding nonrelativistic scheme is maintained in the relativistic domain will be investigated in the following section.

VI. RESULTS

In this section, we test the accuracy of the approximate ROPM scheme, derived within the framework of the \( \text{‘electrostatic case’} \) in the last section, for atomic systems. In order to assess the quality of this approximation, exact results either for the \( xc \) energy \( E_{\text{xc}} \) or for the \( xc \) potential \( V_{xc}(r) \) would be useful. However, for systems where relativistic effects become important, e.g., high-Z atoms, exact results are presently not available. Consequently, we have to look for a different standard to compare with.

Such a standard reference is available within the \( x \)-only limit of RDFT [21,23,24]. As in the nonrelativistic case, the \( x \)-only limit of the \( xc \) energy functional is defined by the use of the exact exchange energy functional, i.e., by the relativistic Fock term, Eq. (6), in the case of only longitudinal (Coulomb) interactions. (Since, in the present context, our principal goal is to test the quality of the RKLI method, we restrict ourselves to this longitudinal case and neglect transverse contributions.) As explained in the preceding section, the exact longitudinal exchange potential \( V_{xc}(r) \) can then be obtained by solving the full ROPM integral equation (15) with \( E_{\text{xc}} \) replaced by \( E_{\text{xc}}^\text{exact} \). Simultaneous solution of the ROPM integral equation and the RKS equation (2) therefore represents the exact implementation of the longitudinal \( x \)-only limit of RDFT. This scheme will serve as a reference standard in the following.

It is first compared — of course — to the RKLI method, which employs the same exact expression (6) for the exchange energy and only approximates the local exchange potential \( V_{xc}^\text{RKLI}(r) \) by means of Eq. (54). Besides that, we list the results from traditional RKS calculations obtained with the longitudinal \( x \)-only RLDA (xRLDA) and two recently introduced relativistic generalized gradient approximation (RGGA) functionals [24]. The first one is based on the Becke88 GGA [42] (RB88), the second one on a GGA functional due to Engel, Chevary, MacDonald, and Vosko [11] (RECMV92).

These various approaches are analyzed for spherical (closed-shell) atoms. To this end, the spin-angular part of the
The RKS wave function is treated analytically and the remaining radial Dirac equation is solved numerically on a logarithmic mesh [21]. In all our calculations we use finite nuclei modeled by a homogeneously charged sphere with radius

\[ R_{\text{mol}} = 1.0793 \, A^{1/3} + 0.735 \, 87 \, \text{fm}, \]

where \( A \) is the atomic mass taken from [43]. We mention in passing that employing finite nuclei is not necessary to ensure convergent results as, for example, in the relativistic Thomas-Fermi model. We incorporate finite nuclei because they represent the physically correct approach.

In Table I, we show the longitudinal ground-state energy \( E^{L}_{\text{tot}} \) obtained from the various self-consistent \( x \)-only RDFT approaches and, in addition, from relativistic Hartree-Fock (RHF) calculations. Comparing the first two columns, we recognize that the RHF and the ROPM data are very close. The largest deviation is found for Be with 41 ppm. With increasing atomic number, the inner orbitals, contributing most to the total energy, become more and more localized such that the difference between the nonlocal RHF potential and the local ROPM decreases. In fact, for No, the difference is down to 2 ppm. We emphasize that these differences are due to the different nature of the two approaches. While the RHF method, by construction, yields the variationally best energy, the ROPM scheme additionally constrains the exchange potential to be local. Consequently, we expect the ROPM results to always be somewhat higher, which is confirmed in Table I. In the third column, the total energies obtained from the RKLI approximation are presented. They always lie above the corresponding ROPM values since the same exchange energy functional is employed in both approaches and the variationally best local potential \( V_{xc}^{L,\text{ROPM}}(r) \) is approximated by Eq. (43) in the RKLI approach. However, the results are clearly seen to agree very closely: For the mean absolute deviation from the exact ROPM data of the 18 neutral atoms listed in Table I, one obtains only 5 mhartree. Thus the RKLI method impressively improves on the RLDA results, for which we find a mean absolute deviation of 6092 mhartree. The accuracy of the RKLI scheme becomes even more obvious when compared to the RGGAs. Both RGGAs improve significantly over the RLDA method. Still, their deviations from the exact ROPM data are more than one order of magnitude larger compared to the RKLI results.

The trends found in the above discussion are almost identical to the ones found in the nonrelativistic case. In order to analyze the relativistic effects more directly, we additionally consider the relativistic contribution to \( E^{L}_{\text{tot}} \), defined by

\[ \Delta E^{L}_{\text{tot}} = E^{L}_{\text{tot}}[n^{R}] - E^{NR}_{\text{tot}}[n^{NR}]. \]

Via this decomposition, we are able to test the quality of the RKLI scheme independently of the accuracy of its nonrelativistic equivalent. Yet, at first, we want to point out that the relativistic treatment leads to drastic corrections especially for high-Z atoms. For example, Table II shows that the relativistic correction of Hg amounts for about 6.7% of the total energy thus demonstrating the need for a fully relativistic treatment. Furthermore, by comparing the second and third columns of Table II, we realize that the ROPM and the RKLI method yield almost identical results for the relativistic contribution \( \Delta E^{L}_{\text{tot}} \). In other words, almost no additional deviations are introduced by the relativistic treatment of the KLI model.
scheme. Considering now the relativistic corrections calculated with the conventional RDFT methods, the conclusions drawn in the discussion of $E_{\text{tot}}^L$ can be repeated: Compared to the RKLI method, the RGGA results are worse by more than one order of magnitude whereas the RLDA is by far least accurate approximation.

These trends also remain valid when other quantities of interest are considered. For example, in Tables III and IV we have listed the relativistic contributions to the longitudinal exchange energy, defined analogously to Eq. (56). Again, the RKLI method almost exactly reproduces the ROPM results. It is worthwhile noting that the exchange energy $E_x^L$ is influenced quite substantially by relativistic effects, too. Taking again Hg as an example, we realize that the 5.8% contribution to $E_x^L$ is of the same order as for the total energy. Furthermore, even for lighter atoms such as Mg, the relativistic corrections to $E_x^L$ are comparable or even larger than the differences between the currently best nonrelativistic exchange functionals. As a consequence, a relativistic treatment is indispensable for the ultimate comparison with experiments [21]. Next, we turn our attention to the calculation of $\Delta E_x$ from the RGGA functionals listed in the third and fourth columns of Table III. With the exception of No, the results are also in excellent agreement with the exact ones. However, when we turn our attention to other systems like some ions of the neon-isoelectronic series shown in Table IV, we recognize that the RGGAs cannot reproduce the exact results to the same level of accuracy as obtained for the neutral atoms above [24]. This can be explained by the fact that the RGGAs are optimized for the neutral atoms. In contrast, the agreement of the RKLI results with the exact ones for these ions is as good as for the neutral atoms.
To further proceed with our investigation, we next raise the question of how well local properties such as the exchange potential $V_x^l(r)$ are reproduced within the different RDFT methods. In Fig. 1, the exchange potential is plotted for the case of Hg. As expected, the RKLI result follows the exact curve most closely. Yet, the pronounced shell structure apparent in the ROPM curve is not fully reproduced, although clearly visible (cf. the more detailed plot in Fig. 2), in the RKLI approximation. However, it again improves significantly over the conventional RDFT results, where the intershell peaks are strongly smeared out or even absent. In Fig. 3, the asymptotic region, which is of particular importance for excitation properties, is plotted in greater detail. There we recover the known deficiencies of the conventional RDFT functionals: The RLDA as well as both RGGAs fall off much too rapidly. In contrast, the RKLI and ROPM curves become indistinguishable in the asymptotic region, both decaying as $1/r$ as $r \to \infty$ and thus reflecting the correct cancellation of self-interaction effects. Since all these observations are already known for the corresponding nonrelativistic functionals, we again consider the relativistic contribution separately. This relativistic contribution to the exchange potential, given by

$$\Delta V_x^l(r) := \frac{V_x^l[n^R](r) - V_x^N[n^R](r)}{V_x^{\text{OPM}}[n^N](r)} \quad ,$$

is plotted in Fig. 4. We first observe strong oscillations between 0.1 a.u. and 5 a.u. These oscillations are introduced by the displacement of the density due to relativistic effects and thus represent a direct consequence of the atomic shell structure. As the shell structure of the exchange potential is not fully reproduced within the RKLI approach, the amplitudes of the oscillations are somewhat smaller compared to $\Delta V_x^{\text{OPM}}(r)$. While these deviations are clearly visible, the RKLI curve is still closest to the exact one, especially in the region near the nucleus and in the asymptotic region, where large deviations occur for the conventional RDFT methods.

The properties of the $x$ potential in the large-$r$ asymptotic region also strongly influence the eigenvalue of the highest occupied orbital shown in Table V: Due to the correct asymptotics, the energies calculated within the RKLI scheme are almost identical to the exact ROPM results. On the contrary, the lack of the correct $1/r$ behavior of the RGGAs and the RLDA shows up in rather poorly reproduced eigenvalues of the highest occupied state. Since in the exact nonrelativistic theory, the highest occupied energy level coincides with the ionization potential [44], we also list this quantity, given by $I - E_{\text{ion}}[N-1] - E_{\text{ion}}[N]$, in Table VI. Again, the RKLI and ROPM methods provide almost identical results. In ad-

![Fig. 1. Longitudinal exchange potential $V_x^l(r)$ for Hg from various self-consistent $x$-only calculations (in hartree units).](image1)

![Fig. 2. Longitudinal exchange potential $V_x^l(r)$ for Hg showing the shell structure of Fig. 1 in more detail (in hartree units).](image2)

![Fig. 3. Asymptotic region of the longitudinal exchange potential of Fig. 1 (in hartree units).](image3)
tion, when comparing the ionization potential to the energies of the highest occupied level, we see that the relation $I - \epsilon_N$ — which we of course expect to hold only approximately, since correlation contributions are neglected — is fulfilled within a few percent for the ROPM and RKLI data, whereas the results of the conventional $x$ functional differ to a much larger extent. To conclude, we note that the inclusion of relativistic effects leads to large corrections for heavy atoms even for the outermost orbitals: Taking Hg as an example, we find a 20% shift for the ionization potential.

TABLE V. Eigenenergy $-\epsilon_N$ of the highest occupied orbital from various self-consistent $x$-only calculations. $\bar{\Delta}$ denotes the mean absolute deviation and $\bar{\delta}$ the average relative deviation (in percent) from the exact ROPM values (all energies are in hartree units).

<table>
<thead>
<tr>
<th></th>
<th>ROPM</th>
<th>OPM</th>
<th>RKLI</th>
<th>RB88</th>
<th>RECMV92</th>
<th>$\bar{\Delta}$</th>
<th>$\bar{\delta}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Be</td>
<td>0.309</td>
<td>0.309</td>
<td>0.309</td>
<td>0.181</td>
<td>0.182</td>
<td>0.170</td>
<td></td>
</tr>
<tr>
<td>Mg</td>
<td>0.253</td>
<td>0.253</td>
<td>0.253</td>
<td>0.149</td>
<td>0.149</td>
<td>0.142</td>
<td></td>
</tr>
<tr>
<td>Ca</td>
<td>0.196</td>
<td>0.196</td>
<td>0.196</td>
<td>0.116</td>
<td>0.116</td>
<td>0.112</td>
<td></td>
</tr>
<tr>
<td>Zn</td>
<td>0.299</td>
<td>0.293</td>
<td>0.298</td>
<td>0.195</td>
<td>0.194</td>
<td>0.191</td>
<td></td>
</tr>
<tr>
<td>Sr</td>
<td>0.181</td>
<td>0.179</td>
<td>0.181</td>
<td>0.108</td>
<td>0.108</td>
<td>0.104</td>
<td></td>
</tr>
<tr>
<td>Cd</td>
<td>0.282</td>
<td>0.266</td>
<td>0.282</td>
<td>0.183</td>
<td>0.182</td>
<td>0.181</td>
<td></td>
</tr>
<tr>
<td>Ba</td>
<td>0.163</td>
<td>0.158</td>
<td>0.163</td>
<td>0.097</td>
<td>0.097</td>
<td>0.095</td>
<td></td>
</tr>
<tr>
<td>Hg</td>
<td>0.329</td>
<td>0.262</td>
<td>0.332</td>
<td>0.223</td>
<td>0.222</td>
<td>0.222</td>
<td></td>
</tr>
</tbody>
</table>

\[
\bar{\Delta} = 0.001 \quad 0.095 \quad 0.095 \quad 0.099
\]

\[
\bar{\delta} = 0.156 \quad 38.3 \quad 38.4 \quad 40.0
\]

VII. CONCLUSIONS

In this work, we derived an approximate ROPM scheme generalizing the arguments of Krieger, Li, and Iafrate to the relativistic domain. As for the full ROPM, the advantage of the RKLI method lies in the fact that xc functionals depending explicitly on a set of RKS single-particle spinors can be treated within the framework of RDFT. This, in particular, allows us to employ the exact expression for the exchange energy functional, i.e., the relativistic Fock term in the longitudinal case. Therefore the RKLI approach satisfies a number of important properties, most notably the freedom from self-interactions implying the correct asymptotic decay.

In numerical tests on spherical atoms within the longitudinal $x$-only limit of RDFT the RKLI method was found to be clearly superior to the known relativistic $x$-only functionals. The results obtained are seen to be in close agreement with the exact ROPM values, thus reducing the deviation of, for example, the widely used RLDA by more than three orders of magnitude. On the other hand, the numerical effort involved is considerably less compared to the solution of the full ROPM scheme. We therefore expect that the RKLI scheme will be successfully used for larger (nonspherical) systems, e.g., molecules, including also correlation contributions.

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