Application of Renormalization to Potential Scattering

Carlos F. de Araujo, Jr.\textsuperscript{a}, Lauro Tomio\textsuperscript{a}, Sadhan K. Adhikari\textsuperscript{a,*}, and T. Frederico\textsuperscript{b}

\textsuperscript{a} Instituto de Física Teórica, Universidade Estadual Paulista, 01405-900 São Paulo, São Paulo, Brasil
\textsuperscript{b} Departamento de Física, Instituto Tecnológico da Aeronáutica, Centro Técnico Aeroespacial, 12228-900 São José dos Campos, São Paulo, Brasil

September 4, 2018

Abstract

A recently proposed renormalization scheme can be used to deal with nonrelativistic potential scattering exhibiting ultraviolet divergence in momentum space. A numerical application of this scheme is made in the case of potential scattering with $r^{-2}$ divergence for small $r$, common in molecular and nuclear physics, by the use of cut-offs in momentum and configuration spaces. The cut-off is finally removed in terms of a physical observable and model-independent result is obtained at low energies. The expected variation of the off-shell behavior of the $t$ matrix arising from the renormalization scheme is also discussed.

PACS Numbers 11.10.Gh, 03.80.+r., 21.30.+y

\textsuperscript{*}John Simon Guggenheim Memorial Foundation Fellow
1 Introduction

Ultraviolet divergences commonly occur in field theories, as a consequence of point-like interactions in the original Lagrangian, both in perturbative expansions and exact solutions. These divergences in perturbative quantum field theory could be eliminated by renormalization to yield a scale \[ \text{(1, 2) and with few exceptions the renormalized perturbative series could not be summed.} \]

In quantum scattering and bound states calculations, a similar situation occurs for potentials with certain singular behavior at short distances. In momentum space such problems develop ultraviolet or large momentum divergence. In these cases, both perturbative and exact methods fail to obtain finite solution.

Motivated by the kind of equations that appear in quantum field theories, the renormalization ideas have been discussed in configuration \[ \text{(3, 4) and momentum \[ \text{(5, 6, 7, 8) spaces for potential scattering, where Dirac delta, contact, or zero-range potentials were used. Such potentials have the advantage to be simple, local and separable, with the additional advantage that all discussion can be analytical \[ \text{(5, 7).} \]

Usually, renormalization is carried out by introducing a sufficiently large (small) cut-off in momentum (configuration) space. With this cut-off the divergences of the original problem are avoided and a meaningful result obtained. In a successful renormalization scheme the knowledge of a physical observable at a particular energy is used to eliminate the cut-off and the final renormalized results should be independent of the detailed renormalization scheme in the limit when the momentum (configuration) cut-off is removed to infinity (zero). Renormalization of these potential models leads to new scales and finite physical observables \[ \text{(3).} \]

The purpose of this work is to perform numerical renormalization of a quantum mechanical problem where analytical solution is not possible. In this way, we can study and understand better the on- and off-shell behavior of the resulting \( t \) matrix, arising from different regularization schemes. We consider renormalization of \( S \)-wave potential scattering with a local attractive potential singular at \( r \to 0 \), given by

\[
V(r) = V_0 \exp\left(-\mu r\right)/r^2.
\]

In addition to performing general numerical renormalization with this potential we shall be particularly interested in studying the expected off-shell behavior of different regularization procedures in the nucleon-nucleon interaction. For a sufficiently large \( |V_0| \), the scattering problem with potential \( (1) \) exhibits ultraviolet divergence and does not permit solutions\[9. \]. At negative energies an infinite number of bound states collapse to infinite binding and at positive energies the scattering equation has noncompact kernel and does not allow scattering solution. For a very small \( |V_0| \) these theoretical problems can be avoided but the problem becomes difficult to handle numerically.

It is appropriate to mention that, in the limit of \( \mu = 0 \), the critical strength given in Ref.\[9] is

\[
V_{0,\text{crit}} = -4 \left( \frac{\hbar^2}{2m} \right) .
\]
where \( m \) is the reduced mass of the two-body system separated by distance \( r \). For attractive (negative) potentials, with \(|V_0| > |V_{0,crit}|\), one encounters an infinite number of bound states with infinite binding. For \(|V_0| < |V_{0,crit}|\), these infinite number of bound states are absent. For a finite non-zero \( \mu \), as the short range divergent behavior of potential (1) is unchanged, similar properties are expected. This behavior is apparent in our numerical results.

The motivation behind studying this potential is that similar potentials occur in various branches of physics, such as molecular, nuclear and atomic physics [10]. Also, from a mathematical point of view, renormalization in quantum mechanics is an interesting field of interest. Potential (1) has particular interest in molecular and nuclear physics. The meson-exchange field-theoretic nucleon-nucleon potentials have terms similar to that in Eq. (1). Potentials with similar behavior also appear in molecular physics. Such terms are treated phenomenologically by introducing an arbitrary cut-off at small \( r \) or large momentum so as to cure the original divergences and to fit some observables. As the problem is mathematically divergent, the results could be very sensitive to the cut-off and it is not clear that if the momentum cut-off can be taken to infinity in order to produce a meaningful renormalization of the problem. In nucleon-nucleon potentials the finite size of the hadrons sets a cut-off in both momentum and configuration spaces, and one should work with a finite cut-off. However, it would be pleasing to see that the results do not get drastically modified as the momentum cut-off is taken to infinity.

In the present work the above problem is renormalized numerically by the introduction of a momentum or configuration space cut-off in different ways. The cut-off is introduced in the Green function or the potential in momentum and configuration spaces, respectively. The cut-off is finally eliminated in favor of a physical observable/information, which is chosen to be the scattering length and the number of bound states of the system. Eventually, we find that by taking the cut-off to infinity, model independent results are obtained. An example of this scheme can be found in Ref.[11], where the binding energy was used as the physical observable. Though in the nucleon-nucleon system one has utmost a single bound state, there could be several bound states in atomic/molecular systems. In this study we have used up to four bound states in order to explore the applicability of the renormalization scheme to these general problems.

The regularization of the solution by modifying the Green function is standard in particle and intermediate energy physics. In atomic, molecular and nuclear physics the traditional method is to modify the potential by form factors. This later procedure is comparable to the modification of the potential at short distances or large momentum as has been done in this work. We demonstrate an equivalence between both schemes for on-shell scattering.

Though most aspects of the present study is quite general, we shall examine some special aspects of interest to nucleon-nucleon scattering. In addition to the on-shell scattering under diverse situations, we also studied the problem of off-shell scattering with potential (1) with a single bound state in some details as that potential simulates nucleon-nucleon scattering. It is well-known that the knowledge of all on- and half-on-shell \( t \) matrix elements is enough to produce the fully off-shell \( t \) matrix elements [12]. This is why we only study the half-shell function [13]. These elements are also interesting from a physical point of view as they are precisely
the elements which can be probed experimentally in a bremsstrahlung experiment [14]. From the present study we determine the extent of the off-shell variation of the phenomenological nucleon-nucleon potentials arising from different regularization schemes.

In this study, in addition to one bound state, we also considered up to four bound states for potential (1). The appearance of few bound states is common in atomic and molecular physics, though in the nucleon-nucleon system commented above there is at most one bound state. The successful renormalization of potential (1) in the presence of several bound states suggests that the present renormalization scheme should be applicable to similar potentials in atomic and molecular physics.

Similar renormalization has been performed recently in quantum mechanics for the one-dimensional $x^{-2}$ potential and the three-dimensional Hulthen potential by Gupta and Rajeev [15], where the interest of such potentials in polymers has been emphasized [16].

The plan of the paper is as follows. In section II, we briefly describe the renormalization scheme in quantum mechanics, in case of a general local potential, following Ref. [7]. In section III we describe the different regularization procedures applied to the scattering equation for the potential given by Eq. (1). Finally, in section IV, we present results and main conclusions.

2 Renormalization in Quantum Mechanics

In operator notation, the Lippmann-Schwinger equation for the $t$ matrix $T$ is given by

$$T(E) = V + VG_0(E)T(E),$$

(3)

where $G_0(E)$ is the free Green function operator,

$$G_0(E) = (E - H_0 + i0)^{-1},$$

(4)

$H_0$ is the free Hamiltonian and $E$ is the center-of-mass energy. $V$ is the potential that, in the present case, is given by Eq. (1).

The partial-wave projection of a matrix element of the operator $O$ ($\equiv T$ or $V$) in momentum space, for a spherically symmetric potential, is given by

$$\langle p|O|q\rangle \equiv \frac{2}{\pi} \sum_{l=0}^{\infty} \frac{4\pi}{2l+1} P_l \left( \frac{p.q}{pq} \right) \sum_{m=-l}^{l} Y_{l,m}^* (\Omega_p) Y_{l,m} (\Omega_q),$$

(5)

where $Y_{l,m}(\Omega_p)$ are the usual spherical harmonics normalized to one, $\Omega_p$ represents the two polar angles of $p$, and $p = |p|$. $P_l(x)$ is the usual Legendre polynomial of order $l$ and $O_l(p,q)$ is the coefficient of the expansion.
Using Eq. (5) in Eq. (3) for a central potential, after the integration in the angular variables (where we use the orthonormalization condition of the spherical harmonics), we obtain the corresponding partial wave projection of Eq. (3):

\[
T_l(p, p'; E) = V_l(p, p') + \frac{2}{\pi} \int_0^\infty q^2 dq V_l(p, q) G_0(q; k^2) T_l(q, p'; E),
\]

where \( E = k^2 \), in units such that \( \hbar = 2m = 1 \), with \( m \) the reduced mass. The momentum-space free Green function is given by

\[
G_0(q; k^2) = (k^2 - q^2 + i0)^{-1}.
\]

With the definition given in Eq. (5), applying the Fourier transform of the local potential \( V(r) \) to momentum space, we can easily obtain the momentum space partial wave coefficients, as detailed in the following. From Eq. (5), by defining \( x \equiv p q / (pq) \) and applying the usual normalization of the Legendre polynomials, we have

\[
V_l(p, q) = \pi^2 \int_{-1}^1 dx P_l(x) \langle p | V | q \rangle = \pi^2 \int_{-1}^1 dx P_l(x) \frac{1}{(2\pi)^3} \int d^3r V(r) \exp(i(p - q) \cdot r) = \frac{1}{2} \int_0^\infty r^2 dr V(r) \int_{-1}^1 dx P_l(x) \frac{\sin(|p - q| r)}{|p - q| r}.
\]

Next, we use the addition theorem given in Eq. (10.1.45) of Ref. [17]:

\[
\frac{\sin(|p - q| r)}{|p - q| r} = j_0 \left( r \sqrt{p^2 + q^2 - 2pqx} \right) = \sum_{n=0}^\infty (2n + 1) j_n(pr) j_n(qr) P_n(x),
\]

where \( j_l(z) \) are spherical Bessel functions of order \( l \) (see chapter 10 of Ref. [17]). Finally, using the normalization condition of the Legendre polynomials, we obtain a simplified expression for our definition of the \( l \)-wave projection of a matrix element of local potential in momentum space:

\[
V_l(p, q) = \int_0^\infty r^2 dr V(r) j_l(pr) j_l(qr).
\]

As we are going to consider only the \( S \)-wave \( (l = 0) \) case, in the following we simplify the notation by neglecting the \( l \) indices.

With the above partial wave projection, the \( S \)-wave scattering length \( a \) and the off-shell function \( f(p, k) \) are defined, respectively, by

\[
a = T(0, 0; 0),
\]

and

\[
f(p, k) = \frac{T(p, k; k^2)}{T(k, k; k^2)}.
\]
In Ref. [7], a constant potential in momentum space, \( V(p', p) = \lambda \), was chosen to exemplify the application of the renormalization scheme in Quantum Mechanics. With such potential the above equation presents ultraviolet divergence, because the kernel of the integral equation is noncompact and Eq. (6) does not have scattering solution.

So, to give meaning to Eq. (6), when it presents ultraviolet divergence, a new regularized equation should be defined, where the Green function or the potential is replaced by a regularized one with a cut-off parameter in momentum space (\( \Lambda \)) or configuration space (\( \alpha \)), respectively. Such regularization truncates the large-momentum parts of this equation and eliminates the ultraviolet divergence. The cut-off parameter is finally eliminated in terms of a physical observable, in order to yield a renormalized \( t \) matrix which should be independent of the regularization procedure. The Green function and potential regularizations should lead to equivalent results, when the cut-off is removed by taking the limit \( \Lambda \rightarrow 0 \) or \( \alpha \rightarrow 0 \), respectively. We show in the following that the phase shifts calculated through a Green function regularization is very similar to those calculated through a potential regularization in the above limit.

One way of introducing a regularized Green function is to multiply the free Green function by a regulator function \( v(q, \Lambda; k) \) that contains a smooth cut-off \( \Lambda(> k) \), such that

\[
G_R(q, \Lambda; k^2) \equiv \frac{v^2(q, \Lambda; k^2)}{k^2 - q^2 + i0}.
\]  

The function \( v(q, \Lambda; k) \) satisfies the conditions

\[
v(k, \Lambda; k) = 1,
\]  

and

\[
\lim_{\Lambda \rightarrow \infty} v(q, \Lambda; k) = 1.
\]  

For an appropriate \( v(q, \Lambda; k) \) the regularized scattering equation,

\[
T_R(p, p'; E) = V(p, p') + \frac{2}{\pi} \int_0^\infty q^2 dq V(p, q) G_R(q, \Lambda; k^2) T_R(q, p'; E),
\]  

has no ultraviolet divergence.

In case of a constant potential in momentum space, the following regulator function was introduced in Ref. [7]:

\[
v^2(q, \Lambda; k) \equiv \frac{(\Lambda^2 + k^2)}{(\Lambda^2 + q^2)}.
\]

For other potentials presenting ultraviolet divergence, one can try the same regulator function \( v(q, \Lambda; k) \), or, (in case of stronger divergence) replace it by another convenient function.

The regularization procedure used above guarantees unitarity because the imaginary part of the Green function is unaffected. However, the limit \( \Lambda \rightarrow \infty \) reduces the regularized Green function to the free Green function and the original ultraviolet divergence reappears. The
renormalized $t$ matrix is obtained by eliminating the cut-off $\Lambda$ in favor of a physical observable at some specific energy.

The above regularization through the Green function, as given in Ref. [7], is on-shell equivalent to a regularization through the potential. To demonstrate this, let us regularize Eq. (6) using the following regularized potential

$$V_R(p, p') = v(p, \Lambda; k) V(p, p') v(p', \Lambda; k). \quad (18)$$

and keeping the Green function unchanged.

Now the regularized $t$-matrix elements satisfy

$$T'_R(p, p'; E) = V_R(p, p') + \frac{2}{\pi} \int_0^\infty q^2 dq V_R(p, q) G_0(q; k^2) T'_R(q, p'; E). \quad (19)$$

The on-shell equivalence of both schemes can be shown by the relation between $T_R$ and $T'_R$, that can be obtained by replacing $V_R$, in Eq. (19), by Eq. (18):

$$T'_R(p, p'; E) = v(p, \Lambda; k) T_R(p, p'; E) v(p', \Lambda; k) \quad (20)$$

$$= v(p, \Lambda; k) \left[ V(p, p') + \frac{2}{\pi} \int_0^\infty q^2 dq V(p, q) \frac{v^2(q, \Lambda; k)}{(k^2 - q^2 + i0)} T_R(q, p'; E) \right] v(p', \Lambda; k).$$

Considering that function $v(p, \Lambda; k)$ should satisfy Eq. (14), all the on-shell scattering observables derived by using $T_R$ or $T'_R$ are identical. Differences between the two approaches can only be detected when calculating off-shell observables.

### 3 Models for Regularization and Renormalization

Fredholm reduction of the Lippmann-Schwinger equation for the $t$ matrix has numerical advantage in dealing with real algebra and one such scheme was used for numerical solution. Auxiliary Fredholm real $\Gamma$-matrix equation was solved numerically [18, 19]. The $\Gamma$-matrix equation avoids the fixed point singularity and has a much weaker kernel. (For such reduced kernel techniques we refer the reader to Refs. [13] and references therein.) When accurate numerical results are needed, the iterative procedure applied to the $\Gamma$ matrix integral equation has considerable advantages, compared to the matrix inversion methods, as shown in Refs. [19].

We apply in this section the renormalization scheme presented in the previous section to solve the scattering equation with potential (8). For a large enough $|V_0|$, this potential has a singular behavior at the origin and admits a ground state with infinite energy, which is physically meaningless.

To test numerically the renormalization scheme we apply four regularization procedures for $S$-wave scattering with this potential. In two of them we have introduced a regulator function in the Green function as described in the previous section. In the other two approaches we introduce a cut-off directly in the configuration space expression of the potential, so that a
simple regulator function in momentum space cannot be factorized. We describe the four schemes (A, B, C, and D) in the following.

**Choice A:** The $S$-wave momentum-space representation of potential (1), after solving Eq. (10) for $l = 0$, with help of Eq. (3.947.2) of Ref. [20], is given by

$$V(p, q) \equiv \frac{V_0}{2pq} \left[ p \arctan \frac{2\mu q}{\mu^2 + p^2 - q^2} + q \arctan \frac{2\mu p}{\mu^2 + q^2 - p^2} + \frac{\mu}{2} \ln \frac{\mu^2 + (p - q)^2}{\mu^2 + (p + q)^2} \right],$$

(21)

with the following limiting behavior

$$V(0, q) = V(q, 0) = \frac{V_0}{q} \arctan \left( \frac{q}{\mu} \right).$$

(22)

The regulator function in the Green function is taken to be the following cut-off step function:

$$v^2(p, \Lambda_A; k) = \Theta(\Lambda_A - p),$$

(23)

where $\Theta(x) = 1$ for $x > 0$ and $=0$ for $x < 0$.

**Choice B:** The potential $V(p, q)$ is also given by Eq. (21), but in this case the regulator function in the Green function is taken to be a smooth cut-off function:

$$v^2(p, \Lambda_B; k) \equiv \frac{(\Lambda_B^2 + k^2)}{(\Lambda_B^2 + p^2)}.$$  

(24)

**Choice C:** The regularization procedure is given in the configuration space, through a cut-off $\alpha_C$, included in the potential, such that in the limit $\alpha_C \to 0$ the regularized potential reduces to Eq. (1). In this case the potential used in calculation is given by

$$V_C(r) = \frac{V_0\exp(-\mu r)}{r(r + \alpha_C)}.$$  

(25)

The corresponding $S$-wave momentum-space matrix element of potential $V_C(p, q)$ is given by

$$V_C(p, q) = \frac{V_0\mu e^{\mu\alpha_C}}{4pq} \int_\mu^\infty dz \exp(-\alpha_C z) \ln \left( \frac{z^2 + (p + q)^2}{z^2 + (p - q)^2} \right),$$

$$= \frac{V_0}{4\alpha_C pq} \ln \left( \frac{\mu^2 + (p + q)^2}{\mu^2 + (p - q)^2} \right) - \frac{2V_0\mu e^{\mu\alpha_C}}{\alpha_C} \int_\mu^\infty dz \frac{z \exp(-\alpha_C z)}{[z^2 + (p + q)^2][z^2 + (p - q)^2]},$$

$$= \frac{V_0\mu e^{\mu\alpha_C}}{4pq} \int_0^1 dy \frac{\exp(-\mu\alpha_C/y)}{y^2} \ln \left( \frac{\mu^2 + (p + q)^2 y^2}{\mu^2 + (p - q)^2 y^2} \right),$$

(26)

with the limiting behavior

$$V_C(p, 0) = V_C(0, p) = V_0\mu e^{\mu\alpha_C} \int_\mu^\infty dz \frac{\exp(-\alpha_C z)}{z^2 + p^2} = V_0\mu e^{\mu\alpha_C} \int_0^1 dy \frac{\exp(-\alpha_C \mu/y)}{\mu^2 + (py)^2}. $$

(27)
Choice D: The procedure for regularization is given in configuration space, through a cut-off \( \alpha_D \), such that in the limit \( \alpha_D \to 0 \), \( V \) reduces to Eq. (1). The regularized potential for this choice is

\[
V_D(r) = \frac{V_0 \exp(-\mu r)}{(r + \alpha_D)^2},
\]

and the corresponding \( S \)-wave momentum-space potential \( V_D(p, q) \) is given by

\[
V_D(p, q) = 2V_0 e^{\mu \alpha_D} \int d\mu \frac{z(z - \mu) \exp(-\alpha_D z)}{[z^2 + (p + q)^2][z^2 + (p - q)^2]} \\
= 2V_0 e^{\mu \alpha_D} \mu^3 \int_0^1 dy \frac{(1 - y) \exp(-\mu \alpha_D / y)}{\mu^4 + 2\mu^2 y^2 (p^2 + q^2) + (p^2 - q^2)^2 y^4}.
\]

For choices A and B the regulator function is introduced in the Green function, as explained in section II. For choices C and D, the cut-off is introduced directly in the configuration space expression of the potential, before the Fourier transform is taken, such that a regulator function cannot be simply factorized in momentum space, as in choices A and B.

The results for the various renormalization schemes is expected to be independent of the detailed regularization procedures. Such expectation is verified in numerical calculation in the next section, and its limitation discussed.

4 Numerical Results and Conclusions

In the following we present the results for phase-shifts, obtained by using four different regularization procedures of the previous section. The reduced mass \( m \) has been taken to be appropriate for the nucleon-nucleon system; \( \hbar^2 / (2m) = 41.47 \text{ MeV} \text{ fm}^2 \). The range parameter \( \mu \) of the potential has been taken to be 1 fm\(^{-1}\). Though we have chosen nuclear units in our calculation, most aspects of this numerical study are of general interest.

In the renormalization scheme we define the cut-off parameter through one of the choices of regularization (for example, choice A), given by Eq. (23), and fix the scattering length to the experimental value (\( a = 5.4 \text{ fm} \), for the \( S \)-wave spin triplet nucleon-nucleon scattering) by adjusting the strength \( V_0 \) of the potential. The scattering length does not fix the number of bound states of the system and the strength of the potential fixes this number. In present renormalization we considered up to four bound states of the system. For other three regularization procedures (choices B, C and D) the strength \( V_0 \) and the corresponding number of bound states (obtained with choice A) are kept fixed together with the scattering length by redefining the corresponding cut-off parameters. With this procedure we guarantee that we are dealing with the renormalization of the same physical interaction via four regularization procedures. Presently, the physical observables that we fix by renormalization are the scattering length and the number of bound states. The numerical parameters used in calculation are given in Table 1.
Though we consider up to four bound states of potential (1) and also consider the limit of very large momentum space cut-off, the situation of interest in the nucleon-nucleon system is the case of one bound state, the deuteron, and a finite cut-off determined by the “size” of the problem. Hence, for the nucleon-nucleon system the first row in Table 1 is of interest, where the momentum space cut-off is set at 10 fm$^{-1}$. We shall explore the off-shell behavior of the nucleon-nucleon scattering under this situation.

In each of Figs. 1-3, we display four sets of curves for phase shift versus energy. In order to have compact figures, in all case we set the zero-energy phase shift at $\pi$. Each set is indicated by an ellipse and labeled by the number of bound states. In each set, there are four curves that represent four different regularization procedures. Schemes (A−D) are conveniently indicated inside the figures. The cut-off parameters of the four schemes are $\Lambda_A$, $\Lambda_B$, $\alpha_C$, and $\alpha_D$. In the renormalization limit the momentum cut-offs $\Lambda_A$ and $\Lambda_B$ should be large, and the configuration cut-offs $\alpha_C$ and $\alpha_D$ should be small. Also, the cut-offs should be large compared to the on-shell momentum. In schemes A − D this means $\Lambda^2_A >> k^2$, $\Lambda^2_B >> k^2$, $\alpha^{-2}_C >> k^2$, and $\alpha^{-2}_D >> k^2$, respectively.

By comparing the curves in each set of Fig. 1, we observe that the phase shifts are numerically very close for energies up to 50 MeV, that correspond to $k^2 \sim 1$ fm$^{-2}$, or $0.01 \Lambda^2_A$. The deviations start to increase as energies approach the cut-off limit.

In Fig. 2, we repeat the same plots as in Fig. 1 with larger momentum cut-offs, $\Lambda_A$ and $\Lambda_B$, or reduced configuration cut-offs, $\alpha_C$ and $\alpha_D$. By increasing the momentum cut-off we approximate the renormalization limit of large cut-off. In this case, comparing the curves in each set, we observe that the phase-shifts are numerically very close for energies up to 500 MeV, that correspond to $k^2 \sim 10$ fm$^{-2} \approx 0.025 \Lambda^2_A$.

The plots of Figs. 1 and 2 are repeated in Fig. 3 for a further increased momentum cut-off (reduced configuration cut-off). As we can see, the phase-shifts, for each set of plots that have the same number of bound states, are very close to each other in all extension of the energies shown in the graph. The energies can go close to 1000 MeV, or $k^2 \sim 20$ fm$^{-2} \approx 0.025\Lambda^2_A$.

As we can see from the three figures, the limit of validity of renormalization increases fast with the increase of the momentum cut-off. As expected, the renormalization become independent of different regularization procedures, if the momentum cut-off is large enough compared to the on shell momentum or equivalently, the configuration cut-off is small enough. In this limit the domain of renormalization extends roughly up to a c.m. energy $E \approx 0.025 E_\Lambda$ where $E_\Lambda = \hbar^2 \Lambda^2/(2m)$.

For $k << \mu$, pure $S$-wave prevails and the scattering phase shift decreases linearly as $k$: $\delta \approx \pi - ka$. This behaviour is consistent with the phase shift of the contact interaction (3): $\tan(\delta) = -ka$, which for small $k$ reduces to $\delta \approx \pi - ka$. This is clearly represented in Figs. 1−3. This behavior does not depend on the details of the potentials and on the number of bound states.

From Table 1, and Figs. 1−3, it is interesting to observe that, even for $\mu \neq 0$, in the case of large cut-offs, the critical strength to produce four bound states is consistent with the limit given in Eq. (2). For the nucleon-nucleon system $V_{0,\text{crit}} = -10.3675$ MeV fm$^2$. Note that in Fig. 3 we have a cut-off with 30 fm$^{-1}$ and the corresponding $V_0 (= -12.482$ MeV.fm$^2$). In the
limit of very large cut-off and large number of bound states, $V_0$ should approach its critical value.

The cut-off behavior of the renormalization scheme is shown in Fig. 4. In this figure we plot the dispersion in phase shifts among the four schemes A − D versus the cut-off parameter $\Lambda$, in the case with four bound states. To show the dispersion in phase shifts we plot the differences between the phase shifts $\Delta \delta$ that show maximum deviation at a fixed energy: choices A and B. The three curves correspond to dispersion at three energies 150 MeV, 500 MeV, and 1000 MeV. As we can observe in this figure, the dispersion approaches zero as we go to infinite momentum cut-off limit, and the results become independent of the regularization procedure that have been used.

We have studied systems that can have up to 4 bound states, but as pointed out before, the situation of interest in the nucleon-nucleon system is the first row in Table 1 where there is only one bound state and a finite momentum space cut-off of about 10 fm$^{-1}$. The situations with more than one bound states are of interest in atomic/molecular physics. Of the three momentum cut-offs used in this work the one at 10 fm$^{-1}$ seems to be more appropriate for the nucleon-nucleon system.

The present work can be used to conclude about the expected off-shell variation of different nucleon-nucleon potential models with the same physical content. Here we have used one constraint on the regularization scheme so as to fit the scattering length. In realistic situations several constraints are used in phenomenological nucleon-nucleon potentials so that the on-shell phase shifts of the different models are practically the same. However, $t$ matrices generated from such potentials may have large off-shell variations as pointed out in Refs. 12, 14. As fully off-shell $t$ matrices are determined in terms of the half-shell $t$ matrices we plot in Fig. 5 the half-shell function (12) for the various schemes at following energies: $E_{c.m.} = 0, 37.5$ MeV, 140 MeV, and 300 MeV. These were the energies explored by Fearing in his study of bremsstrahlung [14] with different realistic potentials.

The half-shell variation in Fig. 5 is not a consequence of the on-shell variations observed in Figs. 1−3. At zero energy the four on-shell $t$ matrices are the same but they may have distinct off-shell behaviors. In the two cases (choices A and B) where momentum cut-offs were employed the regularization schemes are similar and there is no off-shell variation at low energies. The same is true for choices C and D. For the parameters used in the calculation we note that the momentum space cut-off functions (23) and (24) of choices A and B are quite different numerically. The same off-shell behavior generated in these two choices suggest that such momentum space cut-offs in two phenomenological models should not lead to large off-shell variations independent of the values of the parameters used.

The off-shell variation of Fig. 5 is significantly lower than that observed by Fearing in different theoretical models. The larger half-shell variation observed by him was possibly due to differences in the physical contents of the phenomenological nucleon-nucleon potential models. This indicates different intermediate-range (2 − 4 fm) behaviors of those realistic potentials. In the present study the long- and intermediate-range behaviors of the models are kept the same, while the short-range part of these potentials are varied in the different regularization schemes. This variation should increase as the momentum space cut-off $\Lambda$ is reduced. For
a fixed $\Lambda$ it should increase as the momentum variables $k$ or $p$ of the half-shell function are increased. The present half-shell variations noted in Fig. 5 are arising from the inherent differences in different regularization schemes. Similar half-shell variation(s) are expected in different phenomenological nucleon-nucleon potentials with same physical content.

In summary, we have tested numerically renormalization schemes using a singular local potential, that diverges at $r \to 0$ as the inverse of $r^2$. The scattering integral equation exhibits renormalizable ultraviolet divergences in this case. We performed numerical renormalizations with four regularization procedures, two with a cut-off in momentum space and two with a cut-off in configuration space. We find that the renormalization techniques are applicable to this case, such that, for lower energies compared to the cut-off, the results are independent of regularization procedures used to regularize the original equation. In the case of the nucleon-nucleon system the regularization scheme should maintain a single bound state and employ a finite cut-off. In that case we studied the off-shell variation of the different regularization schemes and found that the different phenomenological nucleon-nucleon potential models have larger off-shell variations than found in the present study. Different intermediate-range behaviors of the usual phenomenological potentials are supposed to be responsible for larger off-shell variations. The present renormalization in the presence of up to four bound states is of relevance in atomic and molecular physics, and suggest that similar potentials in atomic and molecular physics can also be successfully renormalized in the present scheme.

We thank the Conselho Nacional de Desenvolvimento Científico e Tecnológico, Fundação de Amparo à Pesquisa do Estado de São Paulo, Coordenação de Aperfeiçoamento de Pessoal do Nível Superior and Financiadora de Estudos e Projetos of Brazil for partial financial support.

References

[1] J. D. Bjorken and S. D. Drell 1965 Relativistic Quantum Fields (New York: McGraw-Hill); R. P. Feynman 1961 Quantum Electrodynamics (New York: Benjamin); C. Itzykson and J. B. Zuber 1980 Quantum Field Theory (New York: McGraw Hill); L. H. Ryder 1985 Quantum Field Theory (Cambridge: Cambridge University Press); P. Ramond 1981 Field Theory: A Modern Primer (Reading: The Benjamin/Cummings Publishing Company)

[2] K. G. Wilson 1971 Phys. Rev. B 4, 3184; K. G. Wilson and J. Kogut 1974 Phys. Rep. 12C, 78

[3] C. Manuel and R. Tarrach 1994 Phys. Lett. B328, 113

[4] P. Gosdzinsky and R. Tarrach 1991 Am. J. Phys. 59, 70; L. R. Mead and J. Godines 1991 ibid. 59, 935

[5] S. K. Adhikari and T. Frederico 1995 Phys. Rev. Lett. 74, 4572

[6] S. Weinberg 1991 Nucl. Phys. B363, 2

[7] S. K. Adhikari, T. Frederico and I. D. Goldman 1995 Phys. Rev. Lett. 74, 487
[8] M. A. B. Beg and R. C. Furlong 1985 Phys. Rev. D 31, 1370
[9] L.D. Landau and E.M. Lifchitz 1977 Quantum Mechanics - Non Relativistic Theory, Chapter V, (New York: Pergamon Press)
[10] W. M. Frank, K. J. Land and R. M. Spector 1971 Rev. Mod. Phys. 43, 1
[11] A.E.A. Amorim, L. Tomio and T. Frederico 1992 Phys. Rev. C 46, 2224
[12] E. F. Redish 1987 Nucl. Phys. A463, 417c
[13] S. K. Adhikari and K. L. Kowalski 1991 Dynamical Collision Theory and its Applications (New York: Academic Press)
[14] H. W. Fearing 1987 Nucl. Phys. A463, 95c
[15] R. Jackiw 1991, in M. A. B. Beg Memorial Volume, Eds. A. Ali and P. Hoodbhoy (Singapore: World Scientific); K. S. Gupta and S. G. Rajeev 1993 Phys. Rev. D 48, 5940; K. M. Case 1950 Phys. Rev. 80, 797; G. Parisi and F. Zirilli 1973 J. Math. Phys. 14, 243
[16] E. Marinari and G. Parisi 1991 Europhys. Lett. 15, 721
[17] M. Abramowitz and I.A. Stegun (eds.) 1968 Handbook of Mathematical Functions (New York: Dover Publications).
[18] L. Tomio and S. K. Adhikari 1980 Phys. Rev. C 22 28; L. Tomio and S. K. Adhikari 1980 ibid C 22 2359
[19] L. Tomio and S. K. Adhikari 1995 Chem. Phys. Lett. 241, 477
[20] I.S. Gradshteyn and I.M. Ryzhik 1965 Table of Integrals, Series, and Products (New York: Academic Press).
Table 1: Numerical parameters used in the four regularization schemes. \( \Lambda_A \) is the regulator used in choice A, \( \Lambda_B \) the corresponding regulator of choice B, \( \alpha_C \) and \( \alpha_D \) are the regulators used in the choices C and D respectively. \( N \) is the number of bound states. For each set of data with the same regulator \( \Lambda_A \), there is a corresponding figure that is indicated in the first column in the left. The scattering length is kept fixed at 5.4 fm in all cases, independent of the number of bound states.

| Fig. | N   | \( V_0 \) (MeV fm\(^2\)) | \( \Lambda_A \) (fm\(^{-1}\)) | \( \Lambda_B \) (fm\(^{-1}\)) | \( \alpha_C \) (fm) | \( \alpha_D \) (fm) |
|------|-----|-----------------------------|-------------------------------|-------------------------------|-------------------|-------------------|
| 1    | 1   | -0.6348                     | 10                            | 12.145                        | 0.020984          | 0.008942          |
| 2    | 2   | -3.6000                     | 10                            | 8.830                         | 0.131200          | 0.044080          |
| 3    | 3   | -14.030                     | 10                            | 6.488                         | 0.403600          | 0.112400          |
| 4    | 4   | -40.132                     | 10                            | 5.200                         | 0.870300          | 0.211380          |
| 1    | 1   | -0.5306                     | 20                            | 25.420                        | 0.008515          | 0.003693          |
| 2    | 2   | -2.2325                     | 20                            | 19.896                        | 0.042190          | 0.015260          |
| 3    | 3   | -7.2221                     | 20                            | 15.395                        | 0.119400          | 0.036880          |
| 4    | 4   | -17.705                     | 20                            | 12.838                        | 0.238200          | 0.065960          |
| 1    | 1   | -0.48935                    | 30                            | 38.930                        | 0.005158          | 0.002261          |
| 3    | 2   | -1.80350                    | 30                            | 31.280                        | 0.022850          | 0.008420          |
| 3    | 3   | -5.39250                    | 30                            | 24.785                        | 0.062700          | 0.020185          |
| 4    | 4   | -12.4820                    | 30                            | 20.917                        | 0.122570          | 0.035700          |

Figure Caption

**Fig. 1** Renormalized phase shifts versus energy using four regularization procedures described in the text. The regulator used for choice A is \( \Lambda_A = 10 \) fm\(^{-1}\). The corresponding regulators for the other three choices, and the corresponding strength \( V_0 \) that maintains the scattering length fixed at 5.4 fm, when varying the number of bound states \( N \), are given in the first set of Table 1.

**Fig. 2** The same as in Fig. 1, using the parameters given in the second set of Table 1, that have \( \Lambda_A = 20 \) fm\(^{-1}\).

**Fig. 3** The same as in Fig. 1, using the parameters given in the third set of Table 1, that have
\[ \Lambda_A = 30 \text{ fm}^{-1}. \]

**Fig. 4** The dispersion in phase shifts \( \Delta \delta \) versus the momentum cut-off \( \Lambda_A \), for the energies 150 MeV, 500 MeV, and 1000 MeV, in the case we have four bound states. We define the dispersion by the difference between the phase shifts of choices A and B, which are enhanced for larger number of bound states \( N \), as shown in Figs. 1-3.

**Fig. 5** Off-shell function \( f(p,k) \) versus \( p \), for different regularization schemes corresponding to the first row of Table 1. The curves are labeled by the respective center of mass energies, in units of MeV.
Phase Shift (radian)

E (MeV)

Choice A - Dashed
Choice B - Solid
Choice C - Dotted
Choice D - Dashed-dotted
Phase Shift (radian) vs. E (MeV)

Choice A - Dashed
Choice B - Solid
Choice C - Dotted
Choice D - Dashed-dotted
Phase Shift (radian)

Choice A - Dashed
Choice B - Solid
Choice C - Dotted
Choice D - Dashed-dotted

E (MeV)
Choices A, B - Solid lines
Choices C, D - Dashed lines