

Basis light-front quantization approach to positronium

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We present the first application of the recently developed basis light-front quantization (BLFQ) method to self-bound systems in quantum field theory, using the positronium system as a test case. Within the BLFQ framework, we develop a two-body effective interaction, operating only in the lowest Fock sector, that implements photon exchange, neglecting fermion self-energy effects. We then solve for the mass spectrum of this interaction at the unphysical coupling $\alpha = 0.3$. The resulting spectrum is in good agreement with the expected Bohr spectrum of nonrelativistic quantum mechanics. We examine in detail the dependence of the results on the regulators of the theory.

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I. INTRODUCTION

The *ab initio* calculation of hadron mass spectra and observables in terms of their underlying quark and gluon degrees of freedom remains a significant challenge to theoretical physics. Bound state problems in a quantum field theory such as QCD are inherently nonperturbative. Hadron problems are also many body in nature and must be solved at strong coupling. Due to the complexity of the problem, seemingly simple quantities, such as the proton magnetic moment, have not yet been accurately calculated.

Observables of interest include the mass spectrum of the hadronic system (in particular “exotic” states with quantum numbers beyond the constituent quark model), along with properties of the corresponding eigenstates. These include transition and decay rates, electric and magnetic moments, form factors, structure functions and generalized parton distributions (GPDs).

Basis light-front quantization (BLFQ) [1] is a promising tool for tackling hadron problems from first principles. BLFQ is a Hamiltonian-based approach that combines the advantages of light-front dynamics [2,3] with modern developments in *ab initio* nuclear structure calculations, such as the no-core shell model (NCSM) [4,5]. The similarity of light-front Hamiltonian quantum field theory to nonrelativistic quantum many-body theory allows the quantum field theoretical bound state problem to be formulated as a large, sparse matrix diagonalization problem. State-of-the-art methods developed for NCSM calculations can then be used to address hadronic systems [6–8]. The diagonalization of the light-front Hamiltonian in a Fock-space basis yields the mass eigenstates of the system, along with amplitudes for evaluating nonperturbative observables.

Recent works have successfully applied BLFQ to the single-electron problem in QED in order to evaluate the electron anomalous magnetic moment both with [9] and without [10,11] an external trap. Another recent application

evaluates the electron GPDs [12]. In addition, BLFQ has been extended to time-dependent strong external field problems such as nonlinear Compton scattering [13,14].

Here, we investigate the positronium system as a test case for applying BLFQ to self-bound systems. Our primary purpose is to confirm that BLFQ is capable of generating the expected Bohr spectrum of positronium, including relativistic effects such as the hyperfine splitting of the ground state.

For this initial test case, we implement a two-body effective interaction that operates only in the lowest Fock sector. The effective interaction implements the one-photon-exchange kernel necessary for Coulomb binding, but neglects the fermion self-energy. In future applications involving dynamical photons in the basis, a nonperturbative renormalization scheme will be needed to deal with fermion self-energy effects. We intend to use the results of this calculation as a benchmark for implementing a Fock-sector dependent renormalization scheme.

We begin by introducing the basic elements of BLFQ, such as our choice of basis and truncation scheme. We then detail the derivation of the two-body effective interaction and solve for the spectrum of positronium using this interaction. Finally, we examine in detail the dependence of the results on the regulators of our theory. Preliminary results were reported in Refs. [15,16].

II. BASIS LIGHT-FRONT QUANTIZATION

In principle, hadron observables can be evaluated by solving the eigenvalue equation

$$P^\mu P_\mu |\Psi\rangle = M^2 |\Psi\rangle, \quad (1)$$

where P^μ is the energy-momentum four-vector operator. In BLFQ, we express the operator P^2 in light-cone gauge. The operator P^2 then plays the role of the Hamiltonian operator in nonrelativistic quantum mechanics. As such, it is sometimes called the “light-cone Hamiltonian” $H_{LC} \equiv P^2$.

(Note that in this convention the Hamiltonian has energy squared units.) This operator can be derived from any field theoretical Lagrangian via the Legendre transform. In BLFQ, Eq. (1) is expressed in a truncated basis, and the resulting finite-dimensional matrix is diagonalized numerically. One then examines the trends in observables as the basis truncation is relaxed to estimate the results in the infinite matrix (or “continuum”) limit.

Of course H_{LC} , being field theoretical in origin, contains terms which change particle number. Thus the basis space for performing a diagonalization must be expanded to include states with any number or species of particles. For example, the positronium wave function could be expressed schematically as

$$|e^+e^- \rangle_{\text{phys}} = a|e^+e^- \rangle + b|e^+e^-\gamma \rangle + c|e^+e^-\gamma\gamma \rangle + d|\gamma \rangle + f|e^+e^-e^+e^- \rangle + \dots \quad (2)$$

When H_{LC} is derived from the QED or QCD Lagrangian, the resulting interactions change particle number by at most 2. The resulting matrix will then be extremely sparse for a many-body calculation.

A. Basis and truncation scheme

In order to numerically diagonalize H_{LC} , the infinite-dimensional basis must be truncated down to a finite dimension. In BLFQ, three separate truncations are made.

First, the number of Fock sectors in the basis is truncated. This truncation will be based on physical as well as practical considerations. For instance, positronium is expected to be fairly well described by the lowest few sectors. Thus, in this introductory work, we limit ourselves to only the $|e^+e^- \rangle$ and $|e^+e^-\gamma \rangle$ sectors. We do not make any attempt here to examine the limit of increasing the number of Fock sectors.

Secondly, we discretize the longitudinal momentum by putting our system in a longitudinal box of length L and applying periodic boundary conditions (BCs). Specifically, we choose periodic BCs for bosons and antiperiodic BCs for fermions. Thus

$$p^+ = \frac{2\pi}{L}j, \quad (3)$$

where j is an integer for bosons, or a half-integer for fermions. For bosons, we exclude the “zero modes,” i.e. $j \neq 0$. In the many-body basis, all basis states are selected to have the same total longitudinal momentum $P^+ = \sum_i p_i^+$, where the sum is over the particles in a particular basis state. We then parametrize P^+ using a dimensionless variable $K = \sum_i j_i$ such that $P^+ = \frac{2\pi}{L}K$. For a given particle i , the longitudinal momentum fraction x is defined as

$$x_i = \frac{p_i^+}{P^+} = \frac{j_i}{K}. \quad (4)$$

Due to the positivity of longitudinal momenta on the light front [3], fixing K also serves as a Fock-space cutoff and makes the number of longitudinal modes finite [17]. It is easy to see that K determines our “resolution” in the longitudinal direction, and thus our resolution on parton distribution functions. The longitudinal continuum limit corresponds to the limit $L, K \rightarrow \infty$.

Finally, in the transverse direction, we employ a 2D harmonic oscillator (HO) basis. On the light front, the generating operator for the 2D HO basis can be expressed as

$$P_+^\Omega = \sum_i \left(\frac{\mathbf{p}_i^2}{2p_i^+} + \frac{\Omega^2}{2} p_i^+ \mathbf{r}_i^2 \right) = \frac{\Omega}{2} \sum_i \left(\frac{\mathbf{p}_i^2}{x_i P^+ \Omega} + x_i P^+ \Omega \mathbf{r}_i^2 \right). \quad (5)$$

Here, and elsewhere, the boldface type is reserved for 2D transverse vectors. Each value of Ω in (5) determines a unique complete and orthonormal basis. By defining the new coordinates [15]

$$\mathbf{q} \equiv \frac{\mathbf{p}}{\sqrt{x}}, \quad \mathbf{s} \equiv \sqrt{x} \mathbf{r}, \quad (6)$$

we can write the generating operator as

$$P_+^\Omega = \frac{\Omega}{2} \sum_i \left[\left(\frac{\mathbf{q}_i}{\sqrt{P^+ \Omega}} \right)^2 + \left(\sqrt{P^+ \Omega} \mathbf{s}_i \right)^2 \right]. \quad (7)$$

We see that P_+^Ω generates a basis of energy scale $b = \sqrt{P^+ \Omega}$. In this work, we use Eq. (7) to define the 2D HO basis states. The momentum-space eigenfunctions of Eq. (7) are

$$\Psi_n^m(\mathbf{q}) = \frac{1}{b} \sqrt{\frac{4\pi \times n!}{(n+|m|)!}} e^{im\phi} \rho^{|m|} e^{-\rho^2/2} L_n^{|m|}(\rho^2), \quad (8)$$

where $\rho \equiv \frac{|\mathbf{q}|}{b}$ and $\phi = \arg(\mathbf{q})$. $L_n^m(x)$ is the generalized (or “associated”) Laguerre polynomial.

The basis is made finite by restricting the number of allowed oscillator quanta in each many-body basis state according to

$$\sum_i (2n_i + |m_i| + 1) \leq N_{\text{max}}. \quad (9)$$

The transverse continuum limit corresponds to $N_{\text{max}} \rightarrow \infty$. In addition, we use an “M-scheme” basis. That is, our many body states have well-defined values of the total angular momentum projection

$$M_J = \sum_i (m_i + s_i), \quad (10)$$

where $s = \pm \frac{1}{2}$ is the fermion spin, but they do not have a well-defined total angular momentum J .

The choice of the 2D HO basis for BLFQ was made with the hadrons in mind. The HO potential is a confining potential, and therefore its wave functions should form an ideal basis for systems subject to QCD confinement. However, it is well known that the HO basis is not ideal for the Coulomb problem, due to the mismatch of the asymptotic behaviors of the wave functions (see, for example, Ref. [18]): the HO wave functions have Gaussian tails, while the hydrogen wave functions have a slower exponential decay. We therefore expect slow convergence in the positronium problem.

B. Center-of-mass factorization

In BLFQ, we construct our many-body basis in single-particle coordinates. The rationale for doing this is its straightforward generalization to a basis of many particles. In principle, relative (Jacobi) coordinates could be used, but this process rapidly becomes intractable as the particle number is increased, due to the need for proper symmetrization of the basis states. Of course, for the two-particle positronium system Jacobi coordinates would be ideal. However, viewing the positronium problem only as a test case for our larger framework, here we nonetheless choose to work in single-particle coordinates.

Since our basis is constructed in single-particle coordinates, the center-of-mass (c.m.) motion of the system is contained in our solutions. The use of the HO basis combined with the N_{\max} truncation is a great advantage here since it allows for the exact factorization of the wave function into ‘‘intrinsic’’ and ‘‘c.m.’’ components, even within a truncated basis. The c.m. motion can then be removed from the low-lying spectrum by introducing a Lagrange multiplier (also known as the Lawson term) to the Hamiltonian [19]. The extra term essentially makes c.m. excitations very costly energetically therefore removing spurious c.m. excitations from the low-lying spectrum.

When the Hamiltonian is expressed in terms of the coordinates (6) exact c.m. factorization is achieved for all eigenstates, even in a basis with arbitrary numbers of sectors, which is the reason for the introduction of these coordinates [15]. The c.m. motion is then governed by

$$H_{\text{c.m.}} = \left(\sum_i \sqrt{x_i} \mathbf{q}_i \right)^2 + b^4 \left(\sum_i \sqrt{x_i} \mathbf{s}_i \right)^2. \quad (11)$$

The c.m. motion can be removed from the low-lying spectrum by adding a Lagrange multiplier proportional to $H_{\text{c.m.}}$ to the Hamiltonian to get

$$H' = H + \lambda(H_{\text{c.m.}} - 2b^2 I), \quad (12)$$

where $H \equiv H_{\text{LC}}$. In practice, one selects λ to be large enough that $2\lambda b^2$ is well above the excitation spectrum of

interest. Demonstrations of the exact c.m. factorization within BLFQ are given in Refs. [15,20].

III. THE HAMILTONIAN

A. Basic structure

We truncate the Fock space to include only $|e^+e^- \rangle$ and $|e^+e^-\gamma \rangle$ states. We wish to formulate an effective potential acting only in the $|e^+e^- \rangle$ space that includes the effects generated by the $|e^+e^-\gamma \rangle$ space. In the formalism of effective potentials, we consider the P space to be the $|e^+e^- \rangle$ space and Q space to be the $|e^+e^-\gamma \rangle$ space. Let \mathcal{P} be the operator that projects onto the P space, and \mathcal{Q} be the operator that projects onto the Q space.

In addition to the effective interaction, the complete Hamiltonian will also include those terms from $H = H_{\text{LC}}$ which act directly in the two-particle space: $\mathcal{P}H\mathcal{P}$. This contains two pieces. First, it contains the two-particle kinetic energy. Secondly, it contains the light-front instantaneous photon-exchange interaction. Thus the total Hamiltonian can be expressed as

$$\langle f | \mathcal{P}H_{\text{tot}}\mathcal{P} | i \rangle = \langle f | \mathcal{P}(H_0 + H_{\text{inst}} + H_{\text{eff}})\mathcal{P} | i \rangle, \quad (13)$$

where states $|i \rangle$ and $|f \rangle$ are states in P space ($|e^+e^- \rangle$). The basis states $|i \rangle$ and $|f \rangle$ are eigenstates of the free Hamiltonian (i.e. $H_0|i \rangle = \epsilon_i|i \rangle$) with eigenvalue

$$\epsilon_i = \sum_j \frac{\mathbf{p}_j^2 + m_j^2}{x_j}, \quad (14)$$

where the sum runs over particles (of mass m_j) in the state $|i \rangle$.

B. Two-body effective interaction

We choose the Bloch form of the effective Hamiltonian. The Bloch Hamiltonian [21] is given by

$$\begin{aligned} \langle f | H_{\text{eff}} | i \rangle &= \frac{1}{2} \sum_n \langle f | \mathcal{P}H\mathcal{Q} | n \rangle \langle n | \mathcal{Q}H\mathcal{P} | i \rangle \\ &\times \left[\frac{1}{\epsilon_i - \epsilon_n} + \frac{1}{\epsilon_f - \epsilon_n} \right]. \end{aligned} \quad (15)$$

Here, $H = H_{\text{LC}} = P^2$ is the light-cone Hamiltonian introduced above. States $|i \rangle$ and $|f \rangle$ are states in P space ($|e^+e^- \rangle$), while state $|n \rangle$ is in the Q space ($|e^+e^-\gamma \rangle$). We have written Eq. (15) in the notation traditionally used for effective potentials in nonrelativistic quantum mechanics calculations. One must remember, though, that both the Hamiltonian H and the ‘‘unperturbed energy’’ ϵ have energy squared units in our formalism. Note that if $|i \rangle = |f \rangle$ this reduces to the usual formula from second-order energy shift in perturbation theory. The derivation of (15), based on a

perturbative expansion of the Okubo-Lee-Suzuki effective Hamiltonian [22–27], is given in Ref. [21].

Since we are interested in primarily the effects of repeated photon exchange, we will only include those combinations of terms in $\mathcal{P}H\mathcal{Q}$ and $\mathcal{Q}H\mathcal{P}$ which generate the photon exchange. We neglect the combinations which result in the photon being emitted and absorbed by the same fermion. That is, we do not incorporate the fermion self-energy. In addition, we work with unit-normalized eigenstates and a fixed value of the coupling constant. Details are presented in the Appendix.

1. Light-front small- x singularities

The instantaneous photon-exchange interaction H_{inst} contains a singularity of the form $\frac{1}{(x_1 - x'_1)^2}$, where x_1 (x'_1) is the longitudinal momentum fraction of the incoming (outgoing) fermion [see Eq. (A19)]. In light-front S -matrix perturbation theory, the amplitude for electron-positron scattering via a dynamical photon contains a term identical to the instantaneous photon-exchange interaction, but with the opposite sign. Thus the instantaneous interaction is canceled *in its entirety*, leaving behind the familiar equal-time Feynman amplitude [2,3]. From this perspective, the instantaneous photon-exchange interaction exists only to cancel this singularity in the light-front electron-positron scattering amplitude. This singularity is an artifact of the use of light-cone gauge, and must be canceled.

In our nonperturbative calculation, this cancellation of small- x singularities does not occur in general. To remove the unphysical singularity in H_{inst} , we introduce a counterterm of the form

$$\langle f|H_{\text{ct}}|i\rangle = -\sum_n \langle f|\mathcal{P}H\mathcal{Q}|n\rangle \langle n|\mathcal{Q}H\mathcal{P}|i\rangle \times \left[\frac{(a-b)^2}{2ab(a+b)} \right], \quad (16)$$

where $a = \epsilon_i - \epsilon_n$ and $b = \epsilon_f - \epsilon_n$. The resulting effective potential, H'_{eff} , is

$$\langle f|H'_{\text{eff}}|i\rangle = \langle f|(H_{\text{eff}} + H_{\text{ct}})|i\rangle = \sum_n \frac{\langle f|\mathcal{P}H\mathcal{Q}|n\rangle \langle n|\mathcal{Q}H\mathcal{P}|i\rangle}{\frac{1}{2}[(\epsilon_i - \epsilon_n) + (\epsilon_f - \epsilon_n)]}. \quad (17)$$

In this form the cancellation of the instantaneous diagram does occur, and $H_{\text{inst}} + H'_{\text{eff}}$ is free of unphysical light-front small- x singularities. Details are given in the Appendix. We note that our choice of counterterm is equivalent to the prescription used in previous work in discretized light-cone quantization (DLCQ) [28–30]. References [28,30] also provide arguments for the plausibility of this prescription.

2. Final expression

Our Hamiltonian is then

$$\langle f|\mathcal{P}H_{\text{tot}}\mathcal{P}|i\rangle = \langle f|\mathcal{P}(H_0 + H_{\text{inst}} + H'_{\text{eff}})\mathcal{P}|i\rangle, \quad (18)$$

where, after canceling the instantaneous interaction, we are left with ($\alpha = g^2/4\pi$)

$$\begin{aligned} H_{\text{inst}} + H'_{\text{eff}} = & \frac{\alpha}{K} \sum_{\bar{\alpha}_1 \bar{\alpha}'_1 \bar{\alpha}_2 \bar{\alpha}'_2} \delta_{j_1+j_2}^{j'_1+j'_2} b_{\bar{\alpha}'_1}^\dagger d_{\bar{\alpha}'_2}^\dagger d_{\bar{\alpha}_2} b_{\bar{\alpha}_1} \\ & \times \sqrt{x_1 x_2 x'_1 x'_2} \int \frac{d^2 \mathbf{q}_1}{(2\pi)^2} \frac{d^2 \mathbf{q}'_1}{(2\pi)^2} \frac{d^2 \mathbf{q}_2}{(2\pi)^2} \frac{d^2 \mathbf{q}'_2}{(2\pi)^2} \\ & \times \frac{(2\pi)^2 \delta^{(2)}(\sqrt{x_1} \mathbf{q}_1 + \sqrt{x_2} \mathbf{q}_2 - \sqrt{x'_1} \mathbf{q}'_1 - \sqrt{x'_2} \mathbf{q}'_2)}{(x_1 - x'_1)^{\frac{1}{2}} [(\epsilon_i - \epsilon_n) + (\epsilon_f - \epsilon_n)]} \\ & \times \Psi_{n_1}^{m_1}(\mathbf{q}_1) \Psi_{n_2}^{m_2}(\mathbf{q}_2) \Psi_{n'_1}^{m'_1*}(\mathbf{q}'_1) \Psi_{n'_2}^{m'_2*}(\mathbf{q}'_2) \\ & \times S_{\alpha_1, \alpha_2, \alpha'_1, \alpha'_2}(\sqrt{x_1} \mathbf{q}_1, \sqrt{x_2} \mathbf{q}_2, \sqrt{x'_1} \mathbf{q}'_1, \sqrt{x'_2} \mathbf{q}'_2), \end{aligned} \quad (19)$$

with $\bar{\alpha}_i$ representing the set of discrete quantum numbers (j_i, s_i, n_i, m_i) and α_i representing the subset (j_i, s_i). If the fermions have mass m_f and the photon has mass μ , the parts of the energy denominator are given by

$$\begin{aligned} \epsilon_i - \epsilon_n &= \frac{(\sqrt{x_1} \mathbf{q}_1)^2 + m_f^2}{x_1} - \frac{(\sqrt{x'_1} \mathbf{q}'_1)^2 + m_f^2}{x'_1} \\ & \quad - \frac{(\sqrt{x_1} \mathbf{q}_1 - \sqrt{x'_1} \mathbf{q}'_1)^2 + \mu^2}{x_1 - x'_1} \\ -(\epsilon_f - \epsilon_n) &= \frac{(\sqrt{x_2} \mathbf{q}_2)^2 + m_f^2}{x_2} - \frac{(\sqrt{x'_2} \mathbf{q}'_2)^2 + m_f^2}{x'_2} \\ & \quad - \frac{(\sqrt{x_2} \mathbf{q}_2 - \sqrt{x'_2} \mathbf{q}'_2)^2 + \mu^2}{x_2 - x'_2}. \end{aligned} \quad (20)$$

The explicit expression for the spinor part, $S_{\alpha_1, \alpha_2, \alpha'_1, \alpha'_2}(\sqrt{x_1} \mathbf{q}_1, \sqrt{x_2} \mathbf{q}_2, \sqrt{x'_1} \mathbf{q}'_1, \sqrt{x'_2} \mathbf{q}'_2)$, is provided in the Appendix (Table I). The highly oscillatory 8D integration can be evaluated using repeated 2D Talmi-Moshinsky transformations [31]. The integral can be reduced down to a single 2D integral, which is evaluated numerically. The details are presented in the Appendix.

Note, also, that a fictitious photon mass μ has been introduced to regulate the expected Coulomb singularity that, while integrable, introduces numerical difficulties. We will later examine the physical limit $\mu \rightarrow 0$.

C. Regulated effective interaction

Previous authors investigating the problem of positronium on the light front with a one-photon-exchange kernel have noted a very small dependence of the ground state

energy on the ultraviolet cutoff of the theory, corresponding to a logarithmic divergence [28,29,32]. References [28,32] both state that the origin of the instability can be traced to a particular term in the effective interaction (or one-photon-exchange kernel) which tends to a nonzero constant, in momentum space, at asymptotically large momentum transfer, corresponding to a Dirac delta potential in coordinate space. Since the 2D Dirac delta potential well has no bound states of finite binding energy [33], this leads to a divergence.

The same divergent term is indeed present in our effective interaction. Numerically, we find that the ground state energy is unstable with increasing N_{\max} , but stable with respect to K for a fixed N_{\max} . The divergence can be removed by introducing a counterterm. Reference [28] argues that such a counterterm can be understood as accounting for the effects of the $|e^+e^-\gamma\gamma\rangle$ Fock sector. The necessary counterterm can be expressed as a modification of the spinor part of the Hamiltonian, $S_{\alpha_1, \alpha_2, \alpha'_1, \alpha'_2}(\sqrt{x_1}\mathbf{q}_1, \sqrt{x_2}\mathbf{q}_2, \sqrt{x'_1}\mathbf{q}'_1, \sqrt{x'_2}\mathbf{q}'_2)$, according to

$$S \rightarrow S - 2[\mathbf{q}_{\text{rel}}^2 + \mathbf{q}'_{\text{rel}}{}^2]\delta_{s_1}^{s'_1}\delta_{s_2}^{s'_2}(\delta_{s_1}^+\delta_{s_2}^- + \delta_{s_1}^-\delta_{s_2}^+), \quad (21)$$

where $s_i = \pm$ represents the fermion spin. In single-particle coordinates, we have $\mathbf{q}_{\text{rel}} = \sqrt{x_2}\mathbf{q}_1 - \sqrt{x_1}\mathbf{q}_2$ and $\mathbf{q}'_{\text{rel}} = \sqrt{x'_2}\mathbf{q}'_1 - \sqrt{x'_1}\mathbf{q}'_2$. We will refer to the resulting interaction as the ‘‘regulated’’ effective interaction. We use the term ‘‘unregulated’’ to refer to the original effective interaction.

IV. NUMERICAL RESULTS

In nonrelativistic quantum mechanics, the hyperfine splitting between the 1S_0 and 3S_1 states of positronium scales as α^4 , where α is the fine structure constant. At physical coupling, the expected hyperfine splitting and even the binding energy are then uncomfortably small relative to the precision of our numerical integrals. Since we would like to use the hyperfine splitting to test our BLFQ results, we use a large coupling of $\alpha = 0.3$ to exaggerate both the binding energy and the hyperfine splitting. We then compare our results not to experiment, but to the predictions of nonrelativistic quantum mechanics at this unphysical value of α . This value of α also allows a direct comparison to the DLCQ results of Refs. [28,29].

The numerical results were obtained using the Hopper Cray XE6 and Edison Cray XC30 at NERSC. ScaLAPACK software [34] was used for the diagonalization. In this particular implementation of BLFQ, the resulting matrix is quite dense. However, in future applications involving multiple Fock sectors, the matrix will be extremely sparse.

We obtain exact c.m. factorization for both the regulated and the unregulated interactions. The spurious c.m. motion is removed by using a Lagrange multiplier, as discussed in Sec. II B, so that all states shown below are in the ground state of c.m. motion.

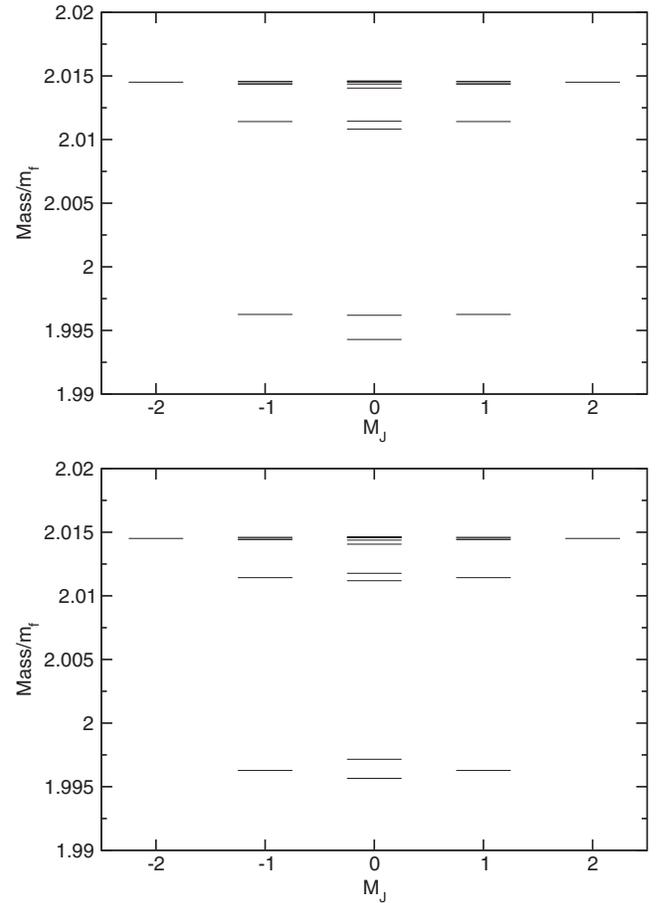


FIG. 1. Representative spectrum of positronium ($\alpha = 0.3$) calculated in BLFQ at $K = N_{\max} = 19$ and $\mu = 0.1m_f$, using the unregulated (top panel) and regulated (bottom panel) effective interactions. The exact energies shown should not be interpreted as the final converged results. Using the unregulated interaction (top panel), the approximate rotational invariance allows for the clear identification of the 1^1S_0 , 1^3S_1 , 2^1S_0 , 2^3S_1 , 2^1P_1 , 2^3P_0 , 2^3P_1 and 2^3P_2 states of the positronium system (see text for details). Using the regulated interaction (bottom panel), the approximate rotational invariance is more strongly broken.

A. Spectrum

A representative spectrum is shown in Fig. 1. These results are produced with $K = N_{\max} = 19$ and $\mu = 0.1m_f$. The energies shown are only representative and should not be considered converged or final. The general features of the spectrum shown here are common to any calculation with $K = N_{\max} = 19$ and above. Convergence will be considered below.

On the light front, the total angular momentum operator is dynamical. In addition, manifest rotational invariance is lost in our calculation due to Fock-sector truncation, as well as the different discretizations in the longitudinal and transverse directions. However, the total J of the states can be extracted by examining the multiplet structure of the spectrum as it appears in the top panel of Fig. 1. The ground

state, for example, appears only in the $M_J = 0$ calculation, suggesting that it has $J = 0$. We also see a triplet of states above the ground state with $M_J = -1, 0, 1$, suggesting that these states form a $J = 1$ multiplet. The lack of manifest rotational invariance is seen only in the lack of exact degeneracy between the states in this multiplet. The difference, however, is quite small, being approximately 1% of the binding energy, and is nearly invisible on this scale. We therefore feel confident extracting J by examining the number of states in each approximately degenerate multiplet.

In the nonrelativistic quantum mechanics description of the positronium system, the states are labeled using the notation $n^{(2S+1)}L_J$, where n is the principal quantum number of atomic physics, S is the total intrinsic spin, L is the total orbital angular momentum (expressed in spectroscopic notation) and J is the total angular momentum. (Note that the n of atomic physics is related to the radial node quantum number of nuclear physics by $n_{\text{atomic}} = n_{\text{radial}} + L + 1$.) The lowest two states are then 1^1S_0 and 1^3S_1 , the singlet and triplet ground state configurations respectively. The splitting between them is referred to as the hyperfine splitting. In a relativistic theory such as BLFQ, L and S are no longer good quantum numbers. We identify the low-lying $J = 0$ and $J = 1$ multiplets in the BLFQ spectrum as the expected 1^1S_0 and 1^3S_1 states of positronium. For simplicity, we will label the states according to their nonrelativistic description, but we stress that we can only extract J (approximately) from our BLFQ calculation, and not L or S . Nevertheless, the BLFQ wave function of the state we identify as the 1^1S_0 state is indeed antisymmetric with respect to spin exchange, as one would expect for an $S = 0$ state. (We reached this conclusion by examining the amplitude by eye at $N_{\text{max}} = 4$.) Similarly, the 1^3S_1 wave function is symmetric with respect to spin exchange, consistent with the nonrelativistic expectation.

A similar grouping of states appear higher up in the spectrum. We identify these as the 2^1S_0 and 2^3S_1 states of positronium. The remaining four states can be identified, via similar reasoning, to be one $J = 0$ multiplet, two $J = 1$ multiplets and one $J = 2$ multiplet. (The highest line in the $M_J = 0$ calculation is thicker to indicate that it represents two nearly degenerate states.) This is exactly what we expect for the $2P$ levels of positronium in nonrelativistic quantum mechanics. We therefore identify these states with the expected 2^1P_1 , 2^3P_0 , 2^3P_1 and 2^3P_2 positronium states. Note that in our BLFQ calculation, we cannot distinguish the two $J = 1$ states because we do not know L and S . (In a relativistic notation, the states are distinguished by their charge conjugation quantum number.)

In the lower panel of Fig. 1, the spectrum of the regulated interaction is shown. The only difference is that the rotational invariance is more severely broken. Compared to the unregulated interaction, the $M_J = 0$ states are shifted upwards, while the $M_J = \pm 1$ states remain essentially

unchanged. A detailed investigation of the rotational symmetry breaking of the two interactions is beyond the scope of this paper. For simplicity, we use the same state identifications for the regulated interaction as for the unregulated interaction.

B. Ground state and hyperfine splitting

We now consider the dependence of the singlet and triplet ground states of positronium on the regulators of our theory. In all of the results below, the value of the basis energy scale is taken to be $b = 0.4m_f$. In principle, the results will converge for any value of b , since the basis is complete regardless of the energy scale. However, in practice, the convergence rate is strongly affected by the value of b . To find the optimal value of the basis energy scale, we plotted the ground state energy as a function of b at $K = N_{\text{max}} = 25$, treating b as a variational parameter. The curve was found to reach a minimum at $b = 0.4m_f$. This value of the basis energy scale, then, is optimal for the convergence of the ground state energy with respect to N_{max} . We emphasize, however, that a different state or observable may have a different optimal value.

Throughout this work we calculate only at *odd* K . The reason is as follows. The kinetic energy term of H_{LC} (14) has a term of the form $m_f^2(\frac{1}{x_1} + \frac{1}{x_2})$. Since we require $x_1 + x_2 = 1$, this term in the kinetic energy takes its minimum value when $x_1 = x_2 = \frac{1}{2}$. This leads to a binding threshold of $2m_f$ as expected. If K is odd, states with $x_1 = x_2$ are present in the basis. If, however, K is even, equal longitudinal momentum splitting is not present in the basis and the ground state energy is increased by a term proportional to m_f^2 . Since, in positronium, the fermion mass scale is much larger than the binding energy scale, the ground state energy is unnecessarily far from convergence if K is even.

In Fig. 2, we plot the singlet and triplet ground state energies as a function $1/N_{\text{max}}$ for a series of values of the longitudinal resolution K and fictitious photon mass μ . The two states are identified in our spectrum as the lowest two states of the $M_J = 0$ calculation, as discussed in Sec. IV A. Since our light-cone Hamiltonian has energy squared units, the energy values plotted are the square root of the eigenvalues of the Hamiltonian. We plot only the results using the regulated interaction, where convergence can be expected. The range of values in $1/N_{\text{max}}$ correspond to $N_{\text{max}} \in [19, 39]$.

The convergence in K is rapid for $\mu \geq 0.6m_f$, as one can see from the fact that the curves for $K = 45$ and $K = 55$ are nearly coincident. As μ is decreased further, the K convergence becomes slower, and we include the $K = 65$ results for comparison. Nevertheless, the results still display a clear converging trend with respect to K for $\mu = 0.02m_f$. The states also exhibit only mild dependence on $1/N_{\text{max}}$, indicating convergence with respect to N_{max} also. To examine the

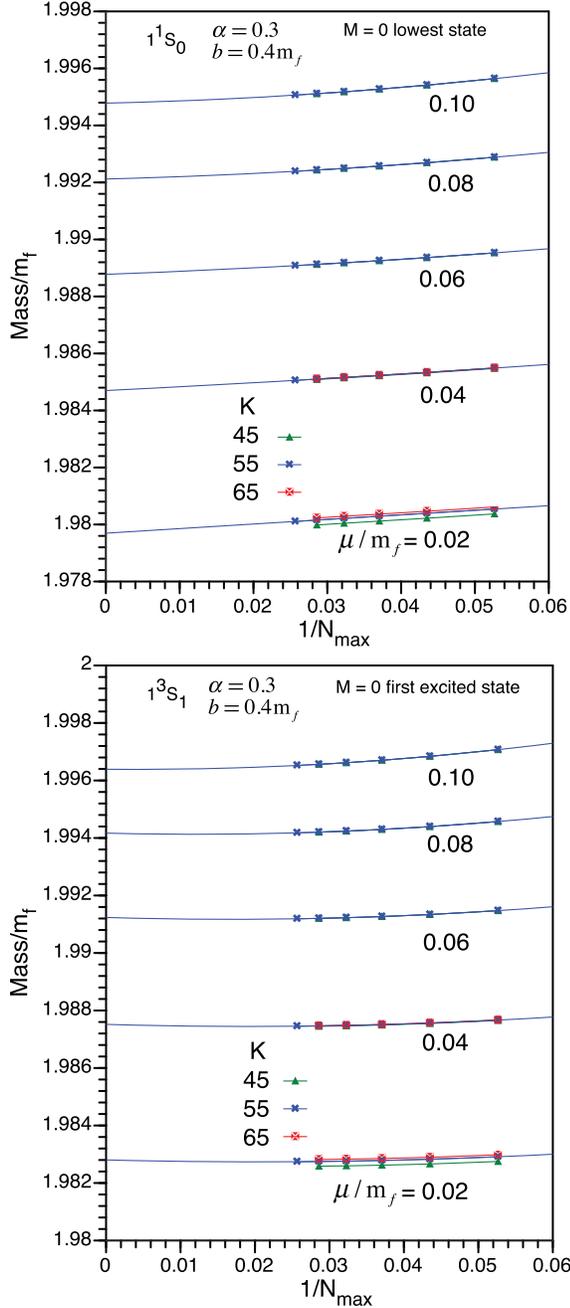


FIG. 2 (color online). Singlet and triplet ground states of positronium ($\alpha = 0.3$) as a function of $1/N_{\max}$ for indicated values of K and μ . The curves are second-order polynomial fits of the $K = 55$ results used to extrapolate the curves to the limit $N_{\max} \rightarrow \infty$ ($1/N_{\max} \rightarrow 0$). The basis energy scale $b = 0.4m_f$ is chosen to optimize the N_{\max} convergence rate for these states.

continuum limit of $N_{\max} \rightarrow \infty$, we fit the $K = 55$ curves to a second-order polynomial in the variable $1/N_{\max}$ and extrapolate the fitted curve to the limit $1/N_{\max} \rightarrow 0$.

The dependence of the states on the fictitious photon mass μ will be discussed in more detail in Sec. IV D, where we examine the physical limit of $\mu \rightarrow 0$ to compare our

results with the standard nonrelativistic quantum mechanics treatment.

C. 2^3P_2 state

As a representative excited state, we also consider the 2^3P_2 state. Being the lowest $J = 2$ state in the positronium spectrum, this state is identified in our spectrum as the lowest state of the $M_J = 2$ calculation. As discussed in Sec. IV B, the optimal value of the HO basis energy scale, b , depends on the specific observable under examination. For the convergence of the ground state energy, we found an optimal value of $b = 0.4m_f$. However, this is not the optimal choice for the 2^3P_2 state. The Coulomb well becomes much wider in position space as one approaches the binding threshold. This indicates that the optimal HO basis length scale for convergence of the 2^3P_2 state energy will be larger than the optimal value for the ground state. A larger length scale, of course, corresponds to a smaller basis energy scale b . Indeed, a variational calculation minimizing the 2^3P_2 state energy, with b as the variational parameter, indicates that $b = 0.1m_f$ would give optimal convergence for this state. We therefore adopt this basis energy scale for our calculations of the 2^3P_2 state binding energy.

The dependence of the 2^3P_2 state on $1/N_{\max}$ for a series of values of K and μ is shown in Fig. 3. The convergence with respect to K is similar to that of the ground state, becoming slower as μ is decreased towards zero. We therefore include higher values of K as μ is decreased to ensure a good estimate of the $K \rightarrow \infty$ limit. The trends in $1/N_{\max}$ are again fit to second-order polynomials to examine the limit $N_{\max} \rightarrow \infty$ ($1/N_{\max} \rightarrow 0$). For $\mu \geq 0.04m_f$, the energy tends to $2m_f$ in the continuum limit, which is the threshold for binding. We conclude that, for $\mu \geq 0.04m_f$, the 2^3P_2 state is a quasibound continuum state, and only becomes bound as the fictitious photon mass decreases below that value.

D. $\mu \rightarrow 0$ limit

In this section, we compare our BLFQ results to the standard formulas of nonrelativistic quantum mechanics. These predictions are (expressed in fermion mass units) [35]

$$M_{\text{singlet}} = 2 - \frac{\alpha^2}{4} \left(1 + \frac{63}{48} \alpha^2 \right) \quad (22)$$

$$M_{\text{triplet}} = 2 - \frac{\alpha^2}{4} \left(1 - \frac{1}{48} \alpha^2 \right) \quad (23)$$

$$M_{3P_2} = 2 - \frac{\alpha^2}{16} \left(1 + \frac{43}{960} \alpha^2 \right). \quad (24)$$

In these formulas, the α^2 term corresponds to the Bohr energies and the α^4 term incorporates first-order perturbative

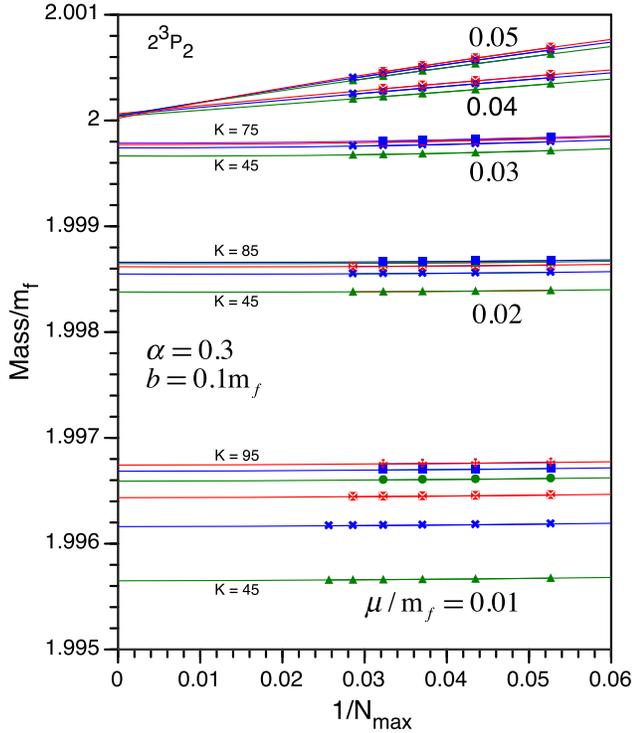


FIG. 3 (color online). The 2^3P_2 state of positronium ($\alpha = 0.3$) as a function of $1/N_{\max}$ for indicated values of K and μ (values of K increase in increments of 10). The discrete points represent the lowest mass eigenstate obtained in calculations with fixed total magnetic projection $M_J = 2$. The curves are second-order polynomial fits used to extrapolate the curves to the limit $N_{\max} \rightarrow \infty$ ($1/N_{\max} \rightarrow 0$). The basis energy scale $b = 0.1m_f$ is chosen to optimize the N_{\max} convergence rate of the 2^3P_2 state binding energy. We include higher values of K as μ is decreased because convergence with K becomes slower. For $\mu \geq 0.04m_f$, the energy converges to $2m_f$ (the threshold for binding) in the continuum limit, indicating that the state is not bound at these high values of μ .

relativistic corrections, neglecting the possibility of a temporary annihilation of the electron and positron into a virtual photon. Since we neglect the $|\gamma\rangle$ Fock sector, our BLFQ results should be comparable to these predictions.

In Fig. 4, we plot the results of our (fixed K) $1/N_{\max} \rightarrow 0$ extrapolations from Secs. IV B and IV C as a function of the fictitious photon mass regulator μ . Recall that the convergence with respect to K gets slower as μ is decreased. To account for this, we increase the fixed value of K used in the $1/N_{\max} \rightarrow 0$ extrapolations as μ gets smaller. For example, we use $K = 55$ for $\mu \geq 0.06m_f$, but $K = 65$ for $\mu = 0.04m_f$ and $\mu = 0.05m_f$. We then use $K = 75$ for $\mu = 0.03m_f$, $K = 85$ for $\mu = 0.02m_f$ and $K = 95$ for $\mu = 0.01m_f$.

We also make a simple extrapolation for the $K \rightarrow \infty$ limit at fixed N_{\max} and fixed μ . This extrapolation was found by comparing many increments in such K values corresponding to vertical sets of symbols in Figs. 2 and 3 at

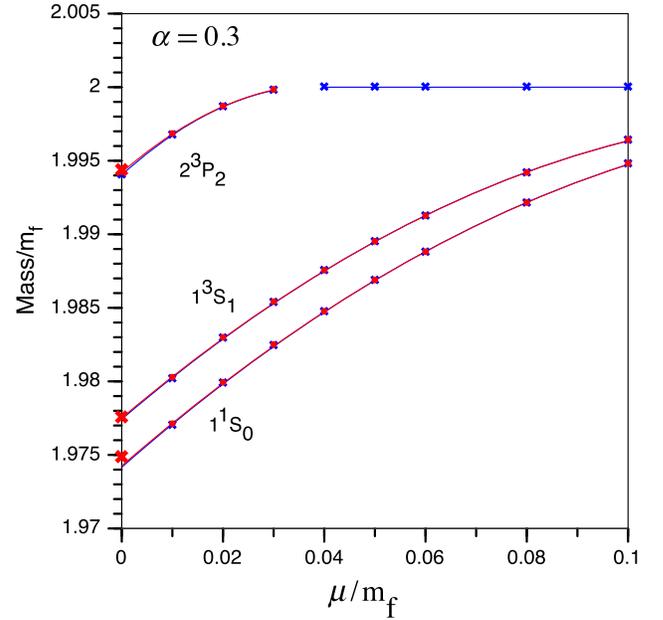


FIG. 4 (color online). Dependence of the 1^1S_0 , 1^3S_1 and 2^3P_2 states on the fictitious photon mass μ . The calculation is performed at a large, unphysical coupling of $\alpha = 0.3$. The plotted blue crosses are the results of the $N_{\max} \rightarrow \infty$ extrapolations (at fixed K) performed in Figs. 2 and 3. The fixed K values chosen for the $1/N_{\max}$ extrapolations are $K = 55$ for $\mu = 0.06\text{--}0.10m_f$, $K = 65$ for $\mu = 0.04\text{--}0.05m_f$, $K = 75$ for $\mu = 0.03m_f$, $K = 85$ for $\mu = 0.02m_f$ and $K = 95$ for $\mu = 0.01m_f$. These values of K are chosen to be sufficiently high such that the results can be considered converged with respect to K . For comparison, the results found using a simple $K \rightarrow \infty$ extrapolation (see text) are shown as solid red circles. The curves are second-order polynomial fits, used to extrapolate to the physical limit $\mu \rightarrow 0$. The crosses on the energy axis represent the predictions of non-relativistic quantum mechanics, including first-order perturbative relativistic corrections, as discussed in the text.

fixed μ values. A stable estimator for the converged result is defined by taking a result at K and adding 1.25 times the difference between that result and the result at $K - 10$. That is, the estimators from several choices of K give a stable result to seven significant figures. The results of these extrapolations are indicated by the solid circles (red dots) in Fig. 4.

We fit the resulting curves to second-order polynomials in μ and extrapolate these curves to the physical limit of $\mu \rightarrow 0$. For comparison, the predictions of Eqs. (22)–(24) are shown as crosses on the vertical axis. The agreement is excellent, despite the large, nonphysical value of $\alpha = 0.3$. The ground state binding energy differs only by $5.84 \times 10^{-4}m_f$, or 2.3% of the binding energy. The hyperfine splitting in BLFQ, $3.26 \times 10^{-3}m_f$, is slightly larger than the expected value of $2.70 \times 10^{-3}m_f$. The 2^3P_2 state binding energy is well reproduced also.

V. CONCLUSIONS AND OUTLOOK

We have presented here the first application of BLFQ to self-bound systems in quantum field theory, using positronium as the test case. In this work, we have truncated the Fock space to include only the $|e^+e^- \rangle$ and $|e^+e^-\gamma \rangle$ sectors. We further restricted the basis to the $|e^+e^- \rangle$ sector alone by developing a two-body effective interaction, incorporating the photon-exchange effects generated by the $|e^+e^-\gamma \rangle$ sector. Diagonalization of the light-cone Hamiltonian including this interaction results in a repeated iteration of the effective interaction, giving a nonperturbative solution to the bound state problem (equivalent to light-front ladder truncation with the one-photon-exchange kernel). In this initial work, we have neglected fermion-self-energy effects arising from the $|e^+e^-\gamma \rangle$ sector.

Positronium is a particularly challenging test case for BLFQ using the HO basis due to the mismatch of the asymptotics of the basis states with the asymptotic bound state amplitudes of Coulombic systems, as discussed in Sec. II A. Nevertheless, we have demonstrated how this mismatch may be overcome with extrapolations to obtain accurate results for the positronium spectrum.

Our converged BLFQ results are in good agreement with the Bohr spectrum of the positronium system at $\alpha = 0.3$. All of the expected total- J multiplets arise in our calculation. In particular, the binding energies (including relativistic corrections) of the 1^1S_0 , 1^3S_1 and 2^3P_2 states are reproduced quantitatively in the BLFQ continuum limit. Thus these results confirm that BLFQ is capable of generating the expected spectrum for positronium. This calculation serves as a strong benchmark test for BLFQ.

A straightforward extension of this work is to include a confining interaction between the fermions. The model should then be applicable to heavy quarkonia systems. A natural choice for the confining potential in BLFQ is a quadratic confinement. Such a quadratic potential is motivated by the phenomenological success of the ‘‘soft wall’’ anti-de Sitter (AdS)/QCD model [36,37]. The effective interaction implemented here would then be interpreted as providing QCD corrections to the semiclassical approximation provided by the AdS/QCD model.

The ultimate goal, however, is to incorporate the higher Fock sectors into the basis, thereby including dynamical bosons. This will be required to obtain realistic QCD results without an *ad hoc* confining interaction. The main challenge is the development and implementation of a nonperturbative renormalization scheme, such as the Fock-sector-dependent scheme of Refs. [38,39]. The recent successful renormalization of the free electron problem in BLFQ [11] with a dynamical photon in the basis suggests the possibility of ‘‘embedding’’ the renormalized, physical electron into the positronium system [40]. Such developments would make BLFQ a powerful tool for addressing the *ab initio* properties of the hadrons.

ACKNOWLEDGMENTS

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APPENDIX: DERIVATION OF TWO-BODY EFFECTIVE INTERACTION

In this appendix, we present the detailed derivation of the two-body effective interaction, including the necessary integrals. Our starting point is the Bloch effective Hamiltonian [21]:

$$\langle f|H_{\text{eff}}|i \rangle = \frac{1}{2} \sum_n \langle f|\mathcal{P}H\mathcal{Q}|n \rangle \langle n|\mathcal{Q}H\mathcal{P}|i \rangle \times \left[\frac{1}{\epsilon_i - \epsilon_n} + \frac{1}{\epsilon_f - \epsilon_n} \right]. \quad (\text{A1})$$

For this BLFQ application, the P space is the $|e^+e^- \rangle$ sector, while the Q space is the $|e^+e^-\gamma \rangle$ sector. Also, in BLFQ applications, $H \equiv H_{\text{LC}} = P^+P^- - \mathbf{P}^2$. Note that \mathbf{P}^2 does not couple the sectors, and so does not contribute to this summation. Therefore, here we will use $H = P^+P^-$. The sum is over the complete Q space. The notation ϵ_i denotes the eigenvalue of the free Hamiltonian H_0 for the state $|i \rangle$ [see Eq. (14)]:

$$\epsilon_i = \sum_j \frac{\mathbf{p}_j^2 + m_j^2}{x_j}, \quad (\text{A2})$$

where the sum runs over all particles (of mass m_j) in the state $|i \rangle$.

After adding the counterterm presented in the main text, the effective Hamiltonian takes the form

$$\langle f|H'_{\text{eff}}|i \rangle = \sum_n \frac{\langle f|\mathcal{P}H\mathcal{Q}|n \rangle \langle n|\mathcal{Q}H\mathcal{P}|i \rangle}{\frac{1}{2}[(\epsilon_i - \epsilon_n) + (\epsilon_f - \epsilon_n)]}. \quad (\text{A3})$$

We will see that this choice, and *only* this choice, results in the cancellation of the light-front small- x divergences.

A. Sum over intermediate states

The basic interaction of the light front quantum electrodynamics Hamiltonian that connects the sectors is the vertex interaction, given by

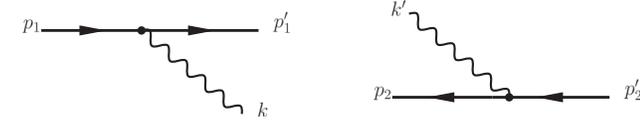


FIG. 5. Basic Hamiltonian interaction vertices $P_{e \to e\gamma}^-$ (left panel) and $P_{\bar{e}\gamma \to \bar{e}}^-$ (right panel).

$$P^- = g \int_{-L}^L dx^- d^2\mathbf{x} \bar{\Psi}(x) \gamma_\mu A^\mu(x) \Psi(x), \quad (\text{A4})$$

evaluated at $x^+ = 0$.

The free field mode expansions in BLFQ are

$$\Psi(x) = \sum_\alpha \frac{1}{\sqrt{2L}} \int \frac{d^2\mathbf{p}}{(2\pi)^2} [b_\alpha(\mathbf{p}) u_s(p) e^{-ipx} + d_\alpha^\dagger(\mathbf{p}) v_s(p) e^{ipx}] \quad (\text{A5})$$

$$A^\mu(x) = \sum_\beta \frac{1}{\sqrt{2Lk^+}} \int \frac{d^2\mathbf{k}}{(2\pi)^2} [a_\beta(\mathbf{k}) \epsilon_\lambda^\mu(k) e^{-ikx} + \text{H.c.}], \quad (\text{A6})$$

where the greek subscripts are shorthand for the set of quantum numbers $\alpha = (j, s)$ and $\beta = (j, \lambda)$, where $s = \pm \frac{1}{2}$ is the fermion spin and where $\lambda = \pm 1$ is the photon helicity.

Plugging these mode expansions into P^- yields terms of the form $P_{e\gamma \to e}^- \sim b^\dagger b a$, $P_{e \to e\gamma}^- \sim b^\dagger b a^\dagger$, $P_{\bar{e}\gamma \to \bar{e}}^- \sim d^\dagger d a$, $P_{\bar{e} \to \bar{e}\gamma}^- \sim d^\dagger d a^\dagger$. These are the only terms which will connect the $|e^+ e^- \rangle$ states to the $|e^+ e^- \gamma \rangle$ states and survive the operation of the projection operators. For example we have the basic interaction vertices shown in Fig. 5, which are given by

$$\begin{aligned} P_{e \to e\gamma}^- &= g \sum_{\alpha_1 \alpha'_1 \beta} \frac{\theta(p_1^+ - p'_1{}^+)}{\sqrt{4\pi K x_k}} \delta_{j_1}^{j_k + j'_1} \\ &\times \int \frac{d^2\mathbf{p}}{(2\pi)^2} \frac{d^2\mathbf{p}'_1}{(2\pi)^2} \frac{d^2\mathbf{k}}{(2\pi)^2} (2\pi)^2 \delta^{(2)}(\mathbf{p}'_1 + \mathbf{k} - \mathbf{p}_1) \\ &\times \bar{u}_{s'_1}(p'_1) \gamma_\mu \epsilon_\lambda^{\mu*}(k) u_{s_1}(p_1) b_{\alpha'_1}^\dagger(\mathbf{p}'_1) a_\beta^\dagger(\mathbf{k}) b_{\alpha_1}(\mathbf{p}_1) \end{aligned} \quad (\text{A7})$$

and

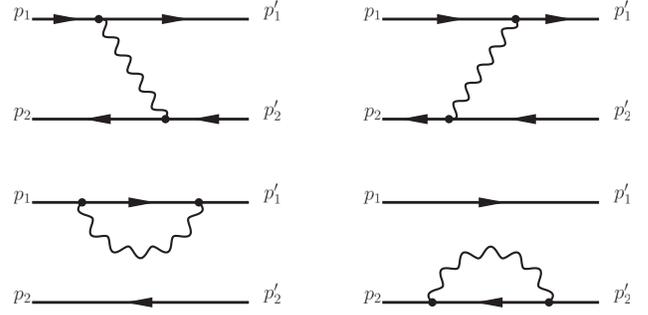


FIG. 6. Iterated interactions generated in the two-body effective interaction. The photon-exchange diagrams correspond to $(H'_{\text{eff}})_{e \to e\gamma}$ (left panel) and $(H'_{\text{eff}})_{\bar{e} \to \bar{e}\gamma}$ (right panel). The fermion-self-energy contributions are neglected in H'_{eff} .

$$\begin{aligned} P_{\bar{e}\gamma \to \bar{e}}^- &= -g \sum_{\alpha_2 \alpha'_2 \beta'} \frac{\theta(p_2'^+ - p_2^+)}{\sqrt{4\pi K x'_k}} \delta_{j_2 + j'_k}^{j'_2} \\ &\times \int \frac{d^2\mathbf{p}_2}{(2\pi)^2} \frac{d^2\mathbf{p}'_2}{(2\pi)^2} \frac{d^2\mathbf{k}'}{(2\pi)^2} (2\pi)^2 \delta^{(2)}(\mathbf{p}_2 + \mathbf{k}' - \mathbf{p}'_2) \\ &\times \bar{v}_{s_2}(p_2) \gamma_\mu \epsilon_\lambda^\mu(k') v_{s'_2}(p'_2) d_{\alpha'_2}^\dagger(\mathbf{p}'_2) d_{\alpha_2}(\mathbf{p}_2) a_{\beta'}(\mathbf{k}'). \end{aligned} \quad (\text{A8})$$

The overall minus in Eq. (A8) has arisen from normal ordering, dropping a vacuum bubble. The theta functions appear since the light-front longitudinal momentum must be positive.

In the effective interaction, these basic interaction vertices are stitched together to generate both fermion-self-energy loops and exchanges of photons between the electron and the positron as shown in Fig. 6. From here on, we neglect these fermion-self-energy terms to focus on the photon-exchange terms. The effective Hamiltonian is then

$$\begin{aligned} H'_{\text{eff}} &= \sum_n \frac{P^+ P_{\bar{e}\gamma \to \bar{e}}^- |n\rangle \langle n| P^+ P_{e \to e\gamma}^-}{\frac{1}{2}[(\epsilon_i - \epsilon_n) + (\epsilon_f - \epsilon_n)]} \\ &+ \sum_n \frac{P^+ P_{e\gamma \to e}^- |n\rangle \langle n| P^+ P_{\bar{e} \to \bar{e}\gamma}^-}{\frac{1}{2}[(\epsilon_i - \epsilon_n) + (\epsilon_f - \epsilon_n)]} \\ &\equiv (H'_{\text{eff}})_{e \to e\gamma} + (H'_{\text{eff}})_{\bar{e} \to \bar{e}\gamma}. \end{aligned} \quad (\text{A9})$$

We now turn to a particular time ordering of the photon exchange, namely the one in which the photon is first emitted by the electron and later absorbed by the positron. This corresponds to the combination $(H'_{\text{eff}})_{e \to e\gamma}$ described above. Let us consider the expression for the effective Hamiltonian in BLFQ. Instead of doing an infinite summation over the complete \mathcal{Q} space in the HO basis, we perform the summation in momentum space, where the sum is an integral and we can write

$$\begin{aligned}
(H'_{\text{eff}})_{e \rightarrow e\gamma} &= \sum_n \frac{P^+ P_{\bar{e}\gamma \rightarrow \bar{e}}^- |n\rangle \langle n| P^+ P_{e \rightarrow e\gamma}^-}{\frac{1}{2}[(\epsilon_i - \epsilon_n) + (\epsilon_f - \epsilon_n)]} \\
&= (P^+)^2 \sum_{\delta \epsilon \xi} \int \frac{d^2 \mathbf{r}}{(2\pi)^2} \frac{d^2 \mathbf{s}}{(2\pi)^2} \frac{d^2 \mathbf{t}}{(2\pi)^2} \frac{P_{\bar{e}\gamma \rightarrow \bar{e}}^- b_{\delta}^{\dagger}(\mathbf{r}) d_{\epsilon}^{\dagger}(\mathbf{s}) a_{\xi}^{\dagger}(\mathbf{t}) |0\rangle \langle 0| a_{\xi}(\mathbf{t}) d_{\epsilon}(\mathbf{s}) b_{\delta}(\mathbf{r}) P_{e \rightarrow e\gamma}^-}{\mathcal{D}_{\delta, \epsilon, \xi}(\mathbf{r}, \mathbf{s}, \mathbf{t})}. \quad (\text{A10})
\end{aligned}$$

The energy denominator is given by

$$\begin{aligned}
\mathcal{D}_{\delta, \epsilon, \xi}(\mathbf{r}, \mathbf{s}, \mathbf{t}) &= \frac{1}{2}[(\epsilon_i - \epsilon_n) + (\epsilon_f - \epsilon_n)] = \frac{\epsilon_i + \epsilon_f}{2} - \epsilon_n \\
&= \frac{1}{2} \left[\left(\frac{\mathbf{p}_1^2 + m_f^2}{x_1} + \frac{\mathbf{p}_2^2 + m_f^2}{x_2} \right) + \left(\frac{\mathbf{p}_1'^2 + m_f^2}{x_1'} + \frac{\mathbf{p}_2'^2 + m_f^2}{x_2'} \right) \right] - \left(\frac{\mathbf{r}^2 + m_f^2}{x_r} + \frac{\mathbf{s}^2 + m_f^2}{x_s} + \frac{\mathbf{t}^2 + \mu^2}{x_t} \right). \quad (\text{A11})
\end{aligned}$$

In the final equality we have used the expression for the kinetic energy appropriate for $H = P^+ P^-$. For simplicity of notation, below we will write $\mathcal{D}_{\delta, \epsilon, \xi}(\mathbf{r}, \mathbf{s}, \mathbf{t}) \equiv \mathcal{D}(r, s, t)$.

We can now substitute Eqs. (A7) and (A8) into Eq. (A10). With the help of the canonical commutation relations

$$[a_{\gamma}(\mathbf{k}), a_{\gamma'}^{\dagger}(\mathbf{k}')]_{-} = (2\pi)^2 \delta^{(2)}(\mathbf{k} - \mathbf{k}') \delta_{\gamma}^{\gamma'} \quad (\text{A12})$$

$$[b_{\alpha}(\mathbf{p}), b_{\alpha'}^{\dagger}(\mathbf{p}')]_{+} = (2\pi)^2 \delta^{(2)}(\mathbf{p} - \mathbf{p}') \delta_{\alpha}^{\alpha'}$$

$$[d_{\alpha}(\mathbf{p}), d_{\alpha'}^{\dagger}(\mathbf{p}')]_{+} = (2\pi)^2 \delta^{(2)}(\mathbf{p} - \mathbf{p}') \delta_{\alpha}^{\alpha'} \quad (\text{A13})$$

it is straightforward to obtain ($\alpha = \frac{\gamma'}{4\pi}$):

$$\begin{aligned}
(H'_{\text{eff}})_{e \rightarrow e\gamma} &= -\frac{\alpha}{K} \sum_{\alpha_1 \alpha_1' \alpha_2 \alpha_2'} \delta_{j_1 + j_2}^{j_1' + j_2'} \theta(p_1^+ - p_1'^+) \int \frac{d^2 \mathbf{p}_1}{(2\pi)^2} \frac{d^2 \mathbf{p}_1'}{(2\pi)^2} \frac{d^2 \mathbf{p}_2}{(2\pi)^2} \frac{d^2 \mathbf{p}_2'}{(2\pi)^2} \frac{(2\pi)^2 \delta^{(2)}(\mathbf{p}_1 + \mathbf{p}_2 - \mathbf{p}_1' - \mathbf{p}_2')}{(x_1 - x_1') \mathcal{D}(p_1', p_2, p_1 - p_1')} \\
&\quad \times (P^+)^2 \bar{u}_{s_1'}(p_1') \gamma_{\mu} u_{s_1}(p_1) \bar{v}_{s_2}(p_2) \gamma_{\nu} v_{s_2'}(p_2') \left(\sum_{\lambda} \epsilon_{\lambda}^{\mu}(k) \epsilon_{\lambda}^{\nu*}(k) \right) b_{\alpha_1'}^{\dagger}(\mathbf{p}_1') d_{\alpha_2'}^{\dagger}(\mathbf{p}_2') |0\rangle \langle 0| d_{\alpha_2}(\mathbf{p}_2) b_{\alpha_1}(\mathbf{p}_1), \quad (\text{A14})
\end{aligned}$$

where $k^{\mu} = (k^-, k^+, \mathbf{k}) \equiv (\frac{\mathbf{p}_1 - \mathbf{p}_1'}{p_1^+ - p_1'^+}, p_1^+ - p_1'^+, \mathbf{p}_1 - \mathbf{p}_1')$. Note that $k^{\mu} \neq (p_1 - p_1')^{\mu}$ due to the minus component. (Light-front energy is not conserved.) Note that in momentum space the state of the photon is completely determined by the external legs, except for the helicity. Thus there remains a sum over helicity states. In the HO basis there would have been an infinite sum over the quantum number n of the oscillator.

Similarly, the other time ordering (corresponding to photon emission by the positron and absorption by the electron) gives

$$\begin{aligned}
(H'_{\text{eff}})_{\bar{e} \rightarrow \bar{e}\gamma} &= -\frac{\alpha}{K} \sum_{\alpha_1 \alpha_1' \alpha_2 \alpha_2'} \delta_{j_1 + j_2}^{j_1' + j_2'} \theta(p_2^+ - p_2'^+) \int \frac{d^2 \mathbf{p}_1}{(2\pi)^2} \frac{d^2 \mathbf{p}_1'}{(2\pi)^2} \frac{d^2 \mathbf{p}_2}{(2\pi)^2} \frac{d^2 \mathbf{p}_2'}{(2\pi)^2} \frac{(2\pi)^2 \delta^{(2)}(\mathbf{p}_1 + \mathbf{p}_2 - \mathbf{p}_1' - \mathbf{p}_2')}{(x_2 - x_2') \mathcal{D}(p_2', p_1, p_2 - p_2')} \\
&\quad \times (P^+)^2 \bar{u}_{s_1'}(p_1') \gamma_{\mu} u_{s_1}(p_1) \bar{v}_{s_2}(p_2) \gamma_{\nu} v_{s_2'}(p_2') \left(\sum_{\lambda} \epsilon_{\lambda}^{\mu}(k) \epsilon_{\lambda}^{\nu*}(k) \right) b_{\alpha_1'}^{\dagger}(\mathbf{p}_1') d_{\alpha_2'}^{\dagger}(\mathbf{p}_2') |0\rangle \langle 0| d_{\alpha_2}(\mathbf{p}_2) b_{\alpha_1}(\mathbf{p}_1). \quad (\text{A15})
\end{aligned}$$

Now, in light-cone gauge, the polarization sum is given by

$$\sum_{\lambda} \epsilon_{\lambda}^{\mu}(k) \epsilon_{\lambda}^{\nu*}(k) = -g^{\mu\nu} + \frac{\eta^{\mu} k^{\nu} + \eta^{\nu} k^{\mu}}{k^+}, \quad (\text{A16})$$

where $\eta^{\mu} = (\eta^-, \eta^+, \boldsymbol{\eta}) = (2, 0, \mathbf{0})$ is a unit vector in the light-front + direction ($k^{\mu} \eta_{\mu} = k^+$). The delta function requires $p_2^+ - p_2'^+ = -(p_1^+ - p_1'^+)$ (similarly for momentum fractions). Also, one can easily show that $\mathcal{D}(p_2', p_1, p_2 - p_2') = -\mathcal{D}(p_1', p_2, p_1 - p_1')$. Therefore, the two time orderings can be combined to obtain our result:

$$\begin{aligned}
H'_{\text{eff}} &= -\frac{\alpha}{K} \sum_{\alpha_1 \alpha_1' \alpha_2 \alpha_2'} \delta_{j_1 + j_2}^{j_1' + j_2'} \int \frac{d^2 \mathbf{p}_1}{(2\pi)^2} \frac{d^2 \mathbf{p}_1'}{(2\pi)^2} \frac{d^2 \mathbf{p}_2}{(2\pi)^2} \frac{d^2 \mathbf{p}_2'}{(2\pi)^2} \frac{(2\pi)^2 \delta^{(2)}(\mathbf{p}_1 + \mathbf{p}_2 - \mathbf{p}_1' - \mathbf{p}_2')}{(x_1 - x_1') \mathcal{D}(p_1', p_2, p_1 - p_1')} \left(-g^{\mu\nu} + \frac{\eta^{\mu} k^{\nu} + \eta^{\nu} k^{\mu}}{k^+} \right) \\
&\quad \times (P^+)^2 \bar{u}_{s_1'}(p_1') \gamma_{\mu} u_{s_1}(p_1) \bar{v}_{s_2}(p_2) \gamma_{\nu} v_{s_2'}(p_2') b_{\alpha_1'}^{\dagger}(\mathbf{p}_1') d_{\alpha_2'}^{\dagger}(\mathbf{p}_2') d_{\alpha_2}(\mathbf{p}_2) b_{\alpha_1}(\mathbf{p}_1). \quad (\text{A17})
\end{aligned}$$

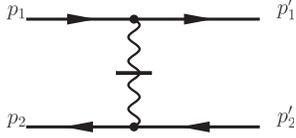


FIG. 7. Instantaneous photon-exchange interaction, P_{inst}^- . The tick mark on the photon denotes instantaneous exchange (in light-front time).

Note that the $|0\rangle\langle 0|$ present in Eqs. (A14) and (A15) is redundant within the $|e^+e^-\rangle$ Fock sector, and has therefore been dropped from Eq. (A17).

B. Canceling the instantaneous photon-exchange term

The instantaneous photon-exchange term in light front quantum electrodynamics is given by

$$P_{\text{inst}}^- = \frac{g^2}{2} \int_{-L}^L dx^- d^2\mathbf{x} \bar{\Psi}(x) \gamma^+ \Psi(x) \frac{1}{(i\partial^+)^2} \bar{\Psi}(x) \gamma^+ \Psi(x), \quad (\text{A18})$$

evaluated at $x^+ = 0$. This is shown diagrammatically in Fig. 7. Again, we substitute in the free field mode expansions and expand. The term of interest is

$$\begin{aligned} H_{\text{inst}} &= P^+ P_{\text{inst}}^- \\ &= -\frac{4\alpha}{K} \sum_{\alpha_1 \alpha_1' \alpha_2 \alpha_2'} \delta_{j_1+j_2}^{j_1'+j_2'} \frac{\delta_{s_1}^{s_1'} \delta_{s_2}^{s_2'}}{(x_1 - x_1')^2} \\ &\quad \times \int \frac{d^2\mathbf{p}_1}{(2\pi)^2} \frac{d^2\mathbf{p}_1'}{(2\pi)^2} \frac{d^2\mathbf{p}_2}{(2\pi)^2} \frac{d^2\mathbf{p}_2'}{(2\pi)^2} (2\pi)^2 \delta^{(2)} \\ &\quad \times (\mathbf{p}_1 + \mathbf{p}_2 - \mathbf{p}_1' - \mathbf{p}_2') \\ &\quad \times b_{\alpha_1}^\dagger(\mathbf{p}_1) d_{\alpha_2}^\dagger(\mathbf{p}_2) d_{\alpha_2}(\mathbf{p}_2) b_{\alpha_1}(\mathbf{p}_1). \end{aligned} \quad (\text{A19})$$

We will now show that the second term in the parentheses of H'_{eff} [Eq. (A17)] exactly cancels H_{inst} . Comparing, we see that cancellation will be achieved if

$$\begin{aligned} A &= \frac{(P^+)^2}{4(x_1 - x_1') \mathcal{D}(p_1', p_2, p_1 - p_1')} \frac{\eta^\mu k^\nu + \eta^\nu k^\mu}{k^+} \\ &\quad \times \bar{u}_{s_1'}(p_1') \gamma_\mu u_{s_1}(p_1) \bar{v}_{s_2}(p_2) \gamma_\nu v_{s_2'}(p_2') = -\frac{\delta_{s_1}^{s_1'} \delta_{s_2}^{s_2'}}{(x_1 - x_1')^2}. \end{aligned} \quad (\text{A20})$$

This is indeed the case, but before we proceed we need a few results. Consider the four-vector $l_e^\mu \equiv (k + p_1' - p_1)^\mu$. By the definition of k^μ [just after Eq. (A14)], we have $l_e^+ = 0$ and $l_e^- = 0$. Thus $l_e^\mu \sim \eta^\mu$. In fact, if we define $Q \equiv l_e^-$, then we have $l_e^\mu = \frac{Q}{2} \eta^\mu$ (recall $\eta^- = 2$). Combining this with the definition of l_e^μ , we see that

$$k^\mu = \frac{Q}{2} \eta^\mu + (p_1 - p_1')^\mu, \quad (\text{A21})$$

where $Q = l_e^- = k^- + p_1'^- - p_1^-$. Similarly, starting from the vector $l_e^\mu \equiv (k + p_2 - p_2')^\mu$, we can show that

$$k^\mu = \frac{\bar{Q}}{2} \eta^\mu + (p_2' - p_2)^\mu, \quad (\text{A22})$$

where $\bar{Q} = l_e^- = k^- + p_2^- - p_2'^-$. Using this information, we see that

$$\begin{aligned} \bar{u}_{s_1'}(p_1') \gamma_\mu k^\mu u_{s_1}(p_1) &= \frac{Q}{2} \bar{u}_{s_1'}(p_1') \gamma^+ u_{s_1}(p_1) = \frac{Q}{2} (2\delta_{s_1}^{s_1'}) \\ \bar{v}_{s_2}(p_2) \gamma_\nu k^\nu v_{s_2'}(p_2') &= \frac{\bar{Q}}{2} \bar{v}_{s_2}(p_2) \gamma^+ v_{s_2'}(p_2') = \frac{\bar{Q}}{2} (2\delta_{s_2}^{s_2'}). \end{aligned} \quad (\text{A23})$$

Use has been made of the Dirac spinor identity $(p - p')^\mu \bar{u}_{s'}(p') \gamma_\mu u_s(p) = 0$ (similarly for the positron) and $\gamma_\mu \eta^\mu = \gamma^+$.

We are now ready to consider

$$\begin{aligned} A &= \frac{(P^+)^2}{4(x_1 - x_1') \mathcal{D}(p_1', p_2, p_1 - p_1')} \frac{\eta^\mu k^\nu + \eta^\nu k^\mu}{k^+} \\ &\quad \times \bar{u}_{s_1'}(p_1') \gamma_\mu u_{s_1}(p_1) \bar{v}_{s_2}(p_2) \gamma_\nu v_{s_2'}(p_2') \\ &= \frac{(2\delta_{s_1}^{s_1'})(2\delta_{s_2}^{s_2'})}{4(x_1 - x_1')^2} \left[\frac{Q + \bar{Q}}{2} P^+ \right] \frac{1}{\mathcal{D}(p_1', p_2, p_1 - p_1')}. \end{aligned} \quad (\text{A24})$$

By writing out the expressions for Q, \bar{Q} and \mathcal{D} , it is readily seen that the factor in brackets is -1 . Therefore

$$A = -\frac{\delta_{s_1}^{s_1'} \delta_{s_2}^{s_2'}}{(x_1 - x_1')^2} \quad (\text{A25})$$

as required for the cancellation of the instantaneous term. Therefore we obtain

$$\begin{aligned} H'_{\text{eff}} + H_{\text{inst}} &= \frac{\alpha}{K} \sum_{\alpha_1 \alpha_1' \alpha_2 \alpha_2'} \delta_{j_1+j_2}^{j_1'+j_2'} \int \frac{d^2\mathbf{p}_1}{(2\pi)^2} \frac{d^2\mathbf{p}_1'}{(2\pi)^2} \frac{d^2\mathbf{p}_2}{(2\pi)^2} \frac{d^2\mathbf{p}_2'}{(2\pi)^2} (2\pi)^2 \delta^{(2)} (\mathbf{p}_1 + \mathbf{p}_2 - \mathbf{p}_1' - \mathbf{p}_2') \\ &\quad \times (P^+)^2 \bar{u}_{s_1'}(p_1') \gamma_\mu u_{s_1}(p_1) \bar{v}_{s_2}(p_2) \gamma^\mu v_{s_2'}(p_2') b_{\alpha_1}^\dagger(\mathbf{p}_1) d_{\alpha_2}^\dagger(\mathbf{p}_2) d_{\alpha_2}(\mathbf{p}_2) b_{\alpha_1}(\mathbf{p}_1), \end{aligned} \quad (\text{A26})$$

where we have contracted the $g^{\mu\nu}$ term.

C. Translation back to harmonic oscillator basis

To translate the result (A26) to the HO basis, we use

$$b_\alpha(\mathbf{p}) = \sum_{nm} b_{\bar{\alpha}} \Psi_n^m(\mathbf{p}), \quad (\text{A27})$$

where $\bar{\alpha}$ is now shorthand for the set of quantum numbers (j, s, n, m) . We therefore obtain the result

$$\begin{aligned} H'_{\text{eff}} + H_{\text{inst}} = & \frac{\alpha}{K} \sum_{\bar{\alpha}_1 \bar{\alpha}'_1 \bar{\alpha}_2 \bar{\alpha}'_2} \delta_{j_1+j_2}^{j'_1+j'_2} b_{\bar{\alpha}'_1}^\dagger b_{\bar{\alpha}'_2}^\dagger d_{\bar{\alpha}_2}^\dagger d_{\bar{\alpha}_1} b_{\bar{\alpha}_1} \int \frac{d^2 \mathbf{p}_1}{(2\pi)^2} \frac{d^2 \mathbf{p}'_1}{(2\pi)^2} \frac{d^2 \mathbf{p}_2}{(2\pi)^2} \frac{d^2 \mathbf{p}'_2}{(2\pi)^2} \frac{(2\pi)^2 \delta^{(2)}(\mathbf{p}_1 + \mathbf{p}_2 - \mathbf{p}'_1 - \mathbf{p}'_2)}{(x_1 - x'_1) \mathcal{D}(p'_1, p_2, p_1 - p'_1)} \\ & \times \Psi_{n_1}^{m_1}(\mathbf{p}_1) \Psi_{n_2}^{m_2}(\mathbf{p}_2) \Psi_{n'_1}^{m'_1}(\mathbf{p}'_1) \Psi_{n'_2}^{m'_2}(\mathbf{p}'_2) (P^+)^2 \bar{u}_{s'_1}(p'_1) \gamma_\mu u_{s_1}(p_1) \bar{v}_{s_2}(p_2) \gamma^\mu v_{s'_2}(p'_2). \end{aligned} \quad (\text{A28})$$

The spinor part $(P^+)^2 \bar{u}_{s'_1}(p'_1) \gamma_\mu u_{s_1}(p_1) \bar{v}_{s_2}(p_2) \gamma^\mu v_{s'_2}(p'_2)$ contains 16 different spin combinations. These are enumerated in Table I.

To translate (A28) into the longitudinal momentum weighted coordinates, we make the substitutions $\mathbf{p} \rightarrow \sqrt{x} \mathbf{q}$, $\int d^2 \mathbf{p} \rightarrow x \int d^2 \mathbf{q}$, $b_{\bar{\alpha}} \rightarrow b_{\bar{\alpha}} / \sqrt{x}$ and $\Psi_n^m(\mathbf{p}) \rightarrow \Psi_n^m(\mathbf{q})$. Equation (A28) now reads

TABLE I. Spinor part $S_{\alpha_1, \alpha_2, \alpha'_1, \alpha'_2}(\mathbf{p}_1, \mathbf{p}_2, \mathbf{p}'_1, \mathbf{p}'_2) \equiv (P^+)^2 \bar{u}_{s'_1}(p'_1) \gamma_\mu u_{s_1}(p_1) \bar{v}_{s_2}(p_2) \gamma^\mu v_{s'_2}(p'_2)$. In this expression, the notation p_j stands for the complex number $p_j = (\mathbf{p}_j)_x + i(\mathbf{p}_j)_y$. Thus $p_j^* = (\mathbf{p}_j)_x - i(\mathbf{p}_j)_y$.

s_1	s_2	s'_1	s'_2	$S_{\alpha_1, \alpha_2, \alpha'_1, \alpha'_2}(\mathbf{p}_1, \mathbf{p}_2, \mathbf{p}'_1, \mathbf{p}'_2)$
+	+	+	+	$2m_f^2 \left(\frac{1}{x_1 x'_1} + \frac{1}{x_2 x'_2} \right) + 2 \left(\frac{p'_1}{x'_1} - \frac{p'_2}{x'_2} \right) \left(\frac{p_1}{x_1} - \frac{p_2}{x_2} \right)$
-	-	-	-	$2m_f^2 \left(\frac{1}{x_1 x'_1} + \frac{1}{x_2 x'_2} \right) + 2 \left(\frac{p_1}{x_1} - \frac{p_2}{x_2} \right) \left(\frac{p'_1}{x'_1} - \frac{p'_2}{x'_2} \right)$
+	-	+	-	$2m_f^2 \left(\frac{1}{x_1 x'_1} + \frac{1}{x_2 x'_2} \right) + 2 \left(\frac{p'_1}{x'_1} - \frac{p_2}{x_2} \right) \left(\frac{p_1}{x_1} - \frac{p'_2}{x'_2} \right)$
-	+	-	+	$2m_f^2 \left(\frac{1}{x_1 x'_1} + \frac{1}{x_2 x'_2} \right) + 2 \left(\frac{p_1}{x_1} - \frac{p'_2}{x'_2} \right) \left(\frac{p'_1}{x'_1} - \frac{p_2}{x_2} \right)$
+	+	+	-	$2m_f \left[\frac{p_2 - p'_2}{x_2 x'_2} + \left(\frac{1}{x_2} - \frac{1}{x'_2} \right) \frac{p_1}{x_1} \right]$
-	+	-	-	$2m_f \left[\frac{p_2 - p'_2}{x_2 x'_2} + \left(\frac{1}{x_2} - \frac{1}{x'_2} \right) \frac{p'_1}{x'_1} \right]$
+	-	+	+	$2m_f \left[\frac{p'_2 - p_2}{x_2 x'_2} + \left(\frac{1}{x'_2} - \frac{1}{x_2} \right) \frac{p'_1}{x'_1} \right]$
-	-	-	+	$2m_f \left[\frac{p'_2 - p_2}{x_2 x'_2} + \left(\frac{1}{x'_2} - \frac{1}{x_2} \right) \frac{p_1}{x_1} \right]$
+	+	-	+	$2m_f \left[\frac{p_1 - p'_1}{x_1 x'_1} + \left(\frac{1}{x_1} - \frac{1}{x'_1} \right) \frac{p_2}{x_2} \right]$
+	-	-	-	$2m_f \left[\frac{p_1 - p'_1}{x_1 x'_1} + \left(\frac{1}{x_1} - \frac{1}{x'_1} \right) \frac{p'_2}{x'_2} \right]$
-	+	+	+	$2m_f \left[\frac{p'_1 - p_1}{x_1 x'_1} + \left(\frac{1}{x'_1} - \frac{1}{x_1} \right) \frac{p'_2}{x'_2} \right]$
-	-	+	-	$2m_f \left[\frac{p'_1 - p_1}{x_1 x'_1} + \left(\frac{1}{x'_1} - \frac{1}{x_1} \right) \frac{p_2}{x_2} \right]$
+	-	-	+	$2m_f^2 \left(\frac{1}{x_1} - \frac{1}{x'_1} \right) \left(\frac{1}{x_2} - \frac{1}{x'_2} \right)$
-	+	+	-	$2m_f^2 \left(\frac{1}{x_1} - \frac{1}{x'_1} \right) \left(\frac{1}{x_2} - \frac{1}{x'_2} \right)$
+	+	-	-	0
-	-	+	+	0

$$\begin{aligned}
H'_{\text{eff}} + H_{\text{inst}} &= \frac{\alpha}{K} \sum_{\bar{\alpha}_1 \bar{\alpha}'_1 \bar{\alpha}_2 \bar{\alpha}'_2} \delta_{j_1+j_2}^{j'_1+j'_2} b_{\bar{\alpha}'_1}^\dagger d_{\bar{\alpha}_1}^\dagger d_{\bar{\alpha}_2}^\dagger b_{\bar{\alpha}'_2} \sqrt{x_1 x_2 x'_1 x'_2} \int \frac{d^2 \mathbf{q}_1}{(2\pi)^2} \frac{d^2 \mathbf{q}'_1}{(2\pi)^2} \frac{d^2 \mathbf{q}_2}{(2\pi)^2} \frac{d^2 \mathbf{q}'_2}{(2\pi)^2} \\
&\times \frac{(2\pi)^2 \delta^{(2)}(\sqrt{x_1} \mathbf{q}_1 + \sqrt{x_2} \mathbf{q}_2 - \sqrt{x'_1} \mathbf{q}'_1 - \sqrt{x'_2} \mathbf{q}'_2)}{(x_1 - x'_1) \mathcal{D}} \Psi_{n_1}^{m_1}(\mathbf{q}_1) \Psi_{n_2}^{m_2}(\mathbf{q}_2) \Psi_{n'_1}^{m'_1}(\mathbf{q}'_1) \Psi_{n'_2}^{m'_2}(\mathbf{q}'_2) \\
&\times S_{\alpha_1, \alpha_2, \alpha'_1, \alpha'_2}(\sqrt{x_1} \mathbf{q}_1, \sqrt{x_2} \mathbf{q}_2, \sqrt{x'_1} \mathbf{q}'_1, \sqrt{x'_2} \mathbf{q}'_2), \tag{A29}
\end{aligned}$$

where the energy denominator factor \mathcal{D} is now given by

$$\mathcal{D} = \frac{1}{2} \left[\left(\frac{(\sqrt{x_1} \mathbf{q}_1)^2 + m_f^2}{x_1} - \frac{(\sqrt{x'_1} \mathbf{q}'_1)^2 + m_f^2}{x'_1} - \frac{(\sqrt{x_1} \mathbf{q}_1 - \sqrt{x'_1} \mathbf{q}'_1)^2 + \mu^2}{x_1 - x'_1} \right) - (1 \rightarrow 2) \right]. \tag{A30}$$

D. The energy denominator integral

First we note that the spinor part simply adds factors of the type $p_i p_j^*$ to the integrations. We can easily absorb these complex transverse momenta into the wave functions using

$$p \Psi_n^m(\mathbf{p}) = b \begin{cases} \sqrt{n+|m|+1} \Psi_n^{m+1}(\mathbf{p}) - \theta(n-1) \sqrt{n} \Psi_{n-1}^{m+1}(\mathbf{p}) & ; m \geq 0 \\ \sqrt{n+|m|} \Psi_n^{m+1}(\mathbf{p}) - \sqrt{n+1} \Psi_{n+1}^{m+1}(\mathbf{p}) & ; m < 0 \end{cases} \tag{A31}$$

$$p^* \Psi_n^m(\mathbf{p}) = b \begin{cases} \sqrt{n+|m|+1} \Psi_n^{m-1}(\mathbf{p}) - \theta(n-1) \sqrt{n} \Psi_{n-1}^{m-1}(\mathbf{p}) & ; m \leq 0 \\ \sqrt{n+|m|} \Psi_n^{m-1}(\mathbf{p}) - \sqrt{n+1} \Psi_{n+1}^{m-1}(\mathbf{p}) & ; m > 0. \end{cases} \tag{A32}$$

Therefore we need only consider the integral

$$\begin{aligned}
I_{ED} &= \int \frac{d^2 \mathbf{q}_1}{(2\pi)^2} \frac{d^2 \mathbf{q}'_1}{(2\pi)^2} \frac{d^2 \mathbf{q}_2}{(2\pi)^2} \frac{d^2 \mathbf{q}'_2}{(2\pi)^2} \frac{(2\pi)^2 \delta^{(2)}(\sqrt{x_1} \mathbf{q}_1 + \sqrt{x_2} \mathbf{q}_2 - \sqrt{x'_1} \mathbf{q}'_1 - \sqrt{x'_2} \mathbf{q}'_2)}{(x_1 - x'_1) \mathcal{D}} \\
&\times \Psi_{n_1}^{m_1}(\mathbf{q}_1) \Psi_{n_2}^{m_2}(\mathbf{q}_2) \Psi_{n'_1}^{m'_1}(\mathbf{q}'_1) \Psi_{n'_2}^{m'_2}(\mathbf{q}'_2). \tag{A33}
\end{aligned}$$

This integral, now expressed in single-particle coordinates, can be simplified by transforming to relative coordinates, using the Talmi-Moshinsky (TM) transform [20,31]:

$$\Psi_{n_1}^{m_1}(\mathbf{q}_1) \Psi_{n_2}^{m_2}(\mathbf{q}_2) = \sum_{NMnm} \mathcal{M}_{n_1 m_1 n_2 m_2}^{NMnm} \Psi_N^M(\mathbf{Q}) \Psi_n^m(\mathbf{q}). \tag{A34}$$

The quantities $\mathcal{M}_{n_1 m_1 n_2 m_2}^{NMnm}$ are known as TM brackets, and the new relative coordinates are

$$\begin{aligned}
\mathbf{q} &= \frac{\sqrt{x_2} \mathbf{q}_1 - \sqrt{x_1} \mathbf{q}_2}{\sqrt{x_1 + x_2}} \\
\mathbf{Q} &= \frac{\sqrt{x_1} \mathbf{q}_1 + \sqrt{x_2} \mathbf{q}_2}{\sqrt{x_1 + x_2}}. \tag{A35}
\end{aligned}$$

We now perform two separate TM transformations from the variables $\mathbf{q}_1, \mathbf{q}_2 \rightarrow \mathbf{Q}, \mathbf{q}$ and $\mathbf{q}'_1, \mathbf{q}'_2 \rightarrow \mathbf{Q}', \mathbf{q}'$. [\mathbf{Q}' and \mathbf{q}' are also given by (A35), with all quantities primed.] In terms of the relative coordinates, the denominator can be expressed as

$$(x_1 - x'_1) \mathcal{D} = -\frac{1}{2(x_1 + x_2)} \left[\left(\sqrt{x'_1 x_2} \mathbf{q} - \sqrt{x_1 x'_2} \mathbf{q}' \right)^2 + \left(\sqrt{x'_1 x_2} \mathbf{q}' - \sqrt{x_1 x'_2} \mathbf{q} \right)^2 + (x_1 + x_2) \Delta \right], \tag{A36}$$

where $\Delta = m_f^2 (x_1 - x'_1)^2 \left[\frac{1}{x_1 x'_1} + \frac{1}{x_2 x'_2} \right] + 2\mu^2 \geq 0$. Note the denominator is independent of \mathbf{Q} and \mathbf{Q}' .

Due to the x dependence of the coordinates \mathbf{q}_i and \mathbf{Q}_i , the TM brackets themselves have x dependence [20]. In particular, the TM phase [20,31] δ is given by $\tan \delta = \sqrt{x_2/x_1}$ for the $\mathbf{q}_1, \mathbf{q}_2 \rightarrow \mathbf{Q}, \mathbf{q}$ transform and by $\tan \delta = \sqrt{x'_2/x'_1}$ for the $\mathbf{q}'_1, \mathbf{q}'_2 \rightarrow \mathbf{Q}', \mathbf{q}'$ transform. The integral then takes the form

$$I_{ED} = \sum_{nmNM} \sum_{n'm'N'M'} \mathcal{M}_{n_1 m_1 n_2 m_2}^{nmNM} \mathcal{M}_{n'_1 (-m'_1) n'_2 (-m'_2)}^{n'm'N'M'} \int \frac{d^2 \mathbf{Q}}{(2\pi)^2} \frac{d^2 \mathbf{q}}{(2\pi)^2} \frac{d^2 \mathbf{Q}'}{(2\pi)^2} \frac{d^2 \mathbf{q}'}{(2\pi)^2} \frac{(2\pi)^2 \delta^{(2)}(\sqrt{x_1+x_2} \mathbf{Q} - \sqrt{x'_1+x'_2} \mathbf{Q}')}{(x_1-x'_1) \mathcal{D}} \times \Psi_n^m(\mathbf{q}) \Psi_N^M(\mathbf{Q}) \Psi_{n'}^{m'}(\mathbf{q}') \Psi_{N'}^{M'}(\mathbf{Q}'). \quad (\text{A37})$$

Since the denominator is independent of \mathbf{Q} and \mathbf{Q}' we can write

$$I_{ED} = \frac{1}{x_1+x_2} \sum_{nmNM} \sum_{n'm'N'M'} \mathcal{M}_{n_1 m_1 n_2 m_2}^{nmNM} \mathcal{M}_{n'_1 (-m'_1) n'_2 (-m'_2)}^{n'm'N'M'} \times \int \frac{d^2 \mathbf{Q}}{(2\pi)^2} \Psi_N^M(\mathbf{Q}) \Psi_{N'}^{M'}(\mathbf{Q}) \int \frac{d^2 \mathbf{q}}{(2\pi)^2} \frac{d^2 \mathbf{q}'}{(2\pi)^2} \frac{\Psi_n^m(\mathbf{q}) \Psi_{n'}^{m'}(\mathbf{q}')}{(x_1-x'_1) \mathcal{D}} \quad (\text{A38})$$

which simplifies to

$$I_{ED} = \frac{1}{x_1+x_2} \sum_{nmNM} \sum_{n'm'N'M'} \mathcal{M}_{n_1 m_1 n_2 m_2}^{nmNM} \mathcal{M}_{n'_1 (-m'_1) n'_2 (-m'_2)}^{n'm'N'M'} \times \delta_N^{N'} \delta_M^{M'} \int \frac{d^2 \mathbf{q}}{(2\pi)^2} \frac{d^2 \mathbf{q}'}{(2\pi)^2} \frac{\Psi_n^m(\mathbf{q}) \Psi_{n'}^{m'}(\mathbf{q}')}{(x_1-x'_1) \mathcal{D}} = \frac{1}{x_1+x_2} \sum_{nmNM} \sum_{n'm'N'M'} \mathcal{M}_{n_1 m_1 n_2 m_2}^{nmNM} \mathcal{M}_{n'_1 (-m'_1) n'_2 (-m'_2)}^{n'm'N(-M)} \times \int \frac{d^2 \mathbf{q}}{(2\pi)^2} \frac{d^2 \mathbf{q}'}{(2\pi)^2} \frac{\Psi_n^m(\mathbf{q}) \Psi_{n'}^{m'}(\mathbf{q}')}{(x_1-x'_1) \mathcal{D}}. \quad (\text{A39})$$

We will refer to the remaining integral as

$$I_{nmn'm'} \equiv \int \frac{d^2 \mathbf{q}}{(2\pi)^2} \frac{d^2 \mathbf{q}'}{(2\pi)^2} \frac{\Psi_n^m(\mathbf{q}) \Psi_{n'}^{m'}(\mathbf{q}')}{(x_1-x'_1) \mathcal{D}}. \quad (\text{A40})$$

$I_{nmn'm'}$ contains terms like $\mathbf{q} \cdot \mathbf{q}'$ in the denominator and therefore has angular dependence. The angular dependence of the denominator can be removed by another TM transform using the variables (here the phase is given by $\tan \delta = 1$)

$$\mathbf{p} = \frac{\mathbf{q} - \mathbf{q}'}{\sqrt{2}} \\ \mathbf{P} = \frac{\mathbf{q} + \mathbf{q}'}{\sqrt{2}}. \quad (\text{A41})$$

The integral is given by

$$I_{nmn'm'} = -2 \sum_{N''M''n''m''} \mathcal{M}_{nmn'm'}^{N''M''n''m''} \times \int \frac{d^2 \mathbf{P}}{(2\pi)^2} \frac{d^2 \mathbf{p}}{(2\pi)^2} \frac{\Psi_{n''}^{m''}(\mathbf{p}) \Psi_{N''}^{M''}(\mathbf{P})}{c\mathbf{P}^2 + c'\mathbf{p}^2 + \Delta}, \quad (\text{A42})$$

where $c = (\sqrt{x'_1 x_2} - \sqrt{x_1 x'_2})^2 / (x_1 + x_2)$ and $c' = (\sqrt{x'_1 x_2} + \sqrt{x_1 x'_2})^2 / (x_1 + x_2)$. If one now substitutes in the expressions for the wave functions, the angular integrations are trivial and it straightforward to obtain

$$I_{nmn'm'} = -\frac{1}{2\pi} \sum_{N''n''} \mathcal{M}_{nmn'm'}^{N''0n''0} \times \int_0^\infty dP dp \frac{e^{-(P+p)/2} L_{N''}(P) L_{n''}(p)}{cP + c'p + \frac{\Delta}{b^2}}, \quad (\text{A43})$$

where $P \equiv \mathbf{P}^2/b^2$ and $p \equiv \mathbf{p}^2/b^2$. $L_n(x) = L_n^0(x)$ are the Laguerre polynomials. The remaining 2D integral cannot be evaluated in closed form and is done numerically. The complete integral, I_{ED} , is then given by

$$I_{ED} = -\frac{1}{2\pi} \frac{1}{x_1+x_2} \sum_{nmNM} \sum_{n'm'N''n''} \mathcal{M}_{n_1 m_1 n_2 m_2}^{nmNM} \times \mathcal{M}_{n'_1 (-m'_1) n'_2 (-m'_2)}^{n'm'N(-M)} \mathcal{M}_{nmn'm'}^{N''0n''0} \times \int_0^\infty dP dp \frac{e^{-(P+p)/2} L_{N''}(P) L_{n''}(p)}{cP + c'p + \frac{\Delta}{b^2}}. \quad (\text{A44})$$

The eightfold sum can be reduced significantly to three folds. Each TM transform comes with two Kronecker deltas:

$$\mathcal{M}_{n_1 m_1 n_2 m_2}^{nmNM} \sim \delta_{2n_1+|m_1|+2n_2+|m_2|}^{2n+|m|+2N+|M|} \delta_{m_1+m_2}^{m+M}. \quad (\text{A45})$$

One of the six Kronecker deltas is left over as an angular momentum conserving the Kronecker delta out front, while the other five can be used to reduce the sum down to three folds. The final result is

$$I_{ED} = -\frac{1}{2\pi} \frac{1}{x_1+x_2} \delta_{m_1+m_2}^{m'_1+m'_2} \sum_{mNN''} \mathcal{M}_{n_1 m_1 n_2 m_2}^{N,\lambda-m,\mu-N,m} \times \mathcal{M}_{n'_1 m'_1 n'_2 m'_2}^{N,\lambda-m,\nu-N,m} \mathcal{M}_{\mu-N,m,\nu-N,-m}^{N'',0,\mu+\nu-2N+|m|-N'',0} \times \int_0^\infty dP dp \frac{e^{-(P+p)/2} L_{N''}(P) L_{\mu+\nu-2N+|m|-N''}(p)}{cP + c'p + \frac{\Delta}{b^2}}, \quad (\text{A46})$$

where $\mu = \mu(m) \equiv n_1 + n_2 + (|m_1| + |m_2| - |m| - |\lambda - m|)/2$, $\nu = \nu(m) \equiv n'_1 + n'_2 + (|m'_1| + |m'_2| - |m| - |\lambda - m|)/2$ and $\lambda = m_1 + m_2$.

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