

Fermionic Mediators, Static Potentials, and Contact/Boundary-Condition Limits

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Abstract

The textbook derivation of a static potential from “field exchange” uses a bosonic mediator linearly sourced by a commuting classical density, yielding an effective action quadratic in the source and (in a static limit) a central Yukawa/Coulomb potential. This derivation does not transplant verbatim to fermionic fields: the linear source terms for fermions require Grassmann-valued sources, so there is no ordinary commuting classical source whose elimination produces a classical potential in the same way. This note isolates the precise obstruction and records the robust infrared replacement: when a microscopic description reduces to local operators at low resolution, the effective interaction is encoded by contact terms (delta kernels and their derivatives) or, equivalently, boundary-condition/self-adjoint-extension data, with renormalization-group running when the contact limit is singular.

This is a dependent note aligned with the broader refinement-compatibility program: contact terms are diagonal-support kernels, and their scale dependence is a compatibility condition rather than an afterthought.

1. Purpose and scope

This note answers a narrowly phrased question: what can it mean for a **fermionic** field to “generate a (central) potential” in the same sense that a massive bosonic field generates a Yukawa potential?

We keep the scope bounded: 1. state the bosonic sourcing \Rightarrow potential mechanism (derivation-first, brief), 2. state the fermionic obstruction precisely (Grassmann sources), 3. give one explicit IR matching witness: **local operators** \Rightarrow **contact/derivative-contact kernels**, 4. connect contact kernels to related point-interaction/RG witnesses.

We do **not** claim that fermions cannot affect forces; we only isolate which parts of the “classical source \Rightarrow potential” story fail, and what the correct replacement statement is at low resolution.

2. What “a field generates a potential” means in the bosonic source story

The archetypal construction is a bosonic mediator φ linearly coupled to a commuting source $J(x)$:

$$S[\varphi; J] = \int d^D x \left(\frac{1}{2} \varphi K \varphi + J \varphi \right), \quad K = (\square + m^2) \text{ (example)}.$$

Integrating out φ (Gaussian elimination) yields an effective action quadratic in the source,

$$S_{\text{eff}}[J] = -\frac{1}{2} \int d^D x d^D y J(x) K^{-1}(x, y) J(y),$$

so the static, nonrelativistic limit of K^{-1} produces a central potential (Yukawa for $m \neq 0$, Coulomb-type when $m = 0$).

The key structural ingredient is that the source is an ordinary commuting function (a classical background density).

Remark 2.1 (Källén-Lehmann spectral representation: the structural backbone of bosonic exchange). The Gaussian-elimination mechanism above has a deeper structural underpinning in the Källén-Lehmann spectral representation (Källén, 1952; Lehmann, 1954). For any scalar field satisfying the Wightman axioms, the full interacting propagator decomposes as $D(p^2) = Z/(p^2 - m_{\text{phys}}^2 + i\epsilon) + \int_{s_0}^{\infty} d\mu^2 \rho(\mu^2)/(p^2 - \mu^2 + i\epsilon)$, where $Z = |\langle 0 | \varphi(0) | 1\text{-particle} \rangle|^2$ is the field-strength renormalization and the spectral function $\rho(\mu^2) \geq 0$ from unitarity (positive-definiteness of the Hilbert-space inner product). For the free field, $\rho = 0$ and $Z = 1$, recovering K^{-1} of Section 2. The static limit of each spectral component gives a Yukawa potential $e^{-\mu r}/r$ with non-negative weight; combined with the negative sign from the Gaussian elimination $S_{\text{eff}} = -(1/2) \int J K^{-1} J$, this yields an attractive static potential for same-sign scalar sources. (For higher-spin exchange, the vertex structure introduces additional sign factors — attractive for even spin, repulsive for odd spin — requiring a separate calculation beyond spectral positivity.) The fermion propagator also admits a spectral representation, but with spinor tensor structure ($\gamma \cdot p$ and scalar components each carrying their own spectral function); the leading spectral function satisfies $\rho_1(\mu^2) \geq 0$ from unitarity, just as in the scalar case. However, the obstruction to “generating a potential” from fermion exchange is not in the spectral weight but in the Grassmann nature of the source terms (Section 3) — the spectral representation exists for both bosons and fermions, but only for bosonic fields does the Gaussian-elimination step produce a classical potential. The spectral representation also underlies the dispersion relation of Remark 3.5: the once-subtracted dispersion integral for $\Pi(q^2)$ is the spectral decomposition of the photon self-energy, with the spectral weight $\text{Im } \Pi(s)$ related to production cross-sections via the optical theorem (Remark 3.6).

Remark 2.2 (Crossing symmetry: the Grassmann obstruction is about the source, not the amplitude). Crossing symmetry relates the same scattering amplitude $\mathcal{A}(s, t, u)$ evaluated in different kinematic regions of the Mandelstam plane — t -channel exchange (spacelike momentum transfer, giving a static potential) and s -channel scattering (timelike, giving production cross-sections) — under the assumption of polynomial-bounded analyticity (the Mandelstam representation, believed to hold in renormalizable QFTs but not proven in full generality). For bosonic exchange, the t -channel amplitude $\mathcal{A}(t) \sim g^2/(t - m^2)$ has a pole whose static limit yields the Yukawa potential of Section 2. The Grassmann obstruction of Section 3 is orthogonal to the channel structure: it prevents the tree-level “classical source elimination” step that reads a c -number potential from the propagator, regardless of whether one works in the s -, t -, or u -channel. The scattering S -matrix for fermion-exchange processes is perfectly well-defined (it involves sums over intermediate states, not classical source elimination), and the loop contributions of Section 3.1 are likewise unaffected. Crossing for amplitudes involving external fermions introduces additional minus signs from Fermi–Dirac statistics and spinor-index structure — the same vertex-structure effects noted in Remark 2.1. The spectral representation of Remark 2.1 provides the unitarity cuts that crossing relates across channels, completing the connection to the dispersion relations of Remark 3.5.

3. Fermionic fields: linear sources are Grassmann, so the classical-source story does not transplant

For a Dirac fermion Ψ , the generating functional with sources is written with **Grassmann-valued** sources $\eta, \bar{\eta}$:

$$Z[\bar{\eta}, \eta] = \int D\bar{\Psi} D\Psi \exp\left(i \int d^D x \bar{\Psi} (i\gamma^\mu \partial_\mu - m)\Psi + i \int d^D x (\bar{\eta} \Psi + \bar{\Psi} \eta)\right).$$

An explicit statement of this form, including that $\eta, \bar{\eta}$ are Grassmann-valued, is recorded in [Floerchinger2024QFT1Lecture21].

Two immediate consequences follow.

Remark 3.1 (Obstruction statement). The bosonic derivation “choose a commuting classical source J , integrate out the field, and read off a classical potential” does not directly apply to fermions, because the linear source terms that couple to Ψ require Grassmann sources rather than commuting c -number densities. Therefore, “fermion exchange generates a classical potential between commuting sources” is not a well-posed transplant of the bosonic story.

This does **not** mean fermions are irrelevant: fermions can and do affect effective interactions through loop effects, through bosonic composite modes (bilinears), and through low-energy EFT operators. The point is that the meaning of “generates a potential” must be stated through one of these controlled mechanisms.

3.1 The controlled alternative: fermion loops modify bosonic propagators

The standard example is vacuum polarization in quantum electrodynamics. A closed electron–positron loop inserted into the photon propagator gives a momentum-dependent correction to the effective electromagnetic coupling,

$$\alpha_{\text{eff}}(q^2) = \frac{\alpha}{1 - \Pi(q^2)}, \quad \Pi(q^2) = -\frac{\alpha}{3\pi} \ln\left(\frac{q^2}{\mu^2}\right) + \dots,$$

where $\Pi(q^2)$ is the vacuum polarization function (the photon self-energy from a fermion one-loop diagram). At low momentum transfer ($|q| \ll m_e$), the loop correction is analytic in q^2 and generates precisely the local operators $C_0 + C_2 q^2 + \dots$ discussed in Section 4 below.

The structural point: fermions affect forces, but the path from “fermion field” to “effective interaction” runs through a quantum loop (not through a tree-level Gaussian elimination of a classical source), and the low-energy residue takes the form of local/contact operators.

Example 3.2 (Uehling potential: the coordinate-space face of vacuum polarization). The momentum-dependent coupling above translates, via Fourier transform, into a coordinate-space correction to the Coulomb potential — the Uehling potential $V_{\text{Uehl}}(r) = -(Z_1 Z_2 \alpha / r) \cdot (2\alpha/3\pi) \int_1^\infty du (1 + 1/(2u^2)) \sqrt{1 - 1/u^2} u^{-1} e^{-2m_e r u}$. At short distances ($r \ll 1/m_e$) the integral yields a logarithmic correction $\propto \ln(1/(m_e r))$, reflecting the running coupling and matching the analytic $C_0 + C_2 q^2 + \dots$ expansion of Section 4; at long distances ($r \gg 1/m_e$) it is exponentially suppressed $\propto e^{-2m_e r}$, confirming that the fermion decouples below its mass threshold. The dominant observable consequence is the vacuum-polarization contribution to the hydrogen Lamb shift (≈ 27 MHz of the total ≈ 1058 MHz $2S$ – $2P$ splitting), which probes the modified potential at nuclear distances where the S -wave wavefunction satisfies $|\psi(0)|^2 \neq 0$.

Remark 3.3 (Anomalous magnetic moment: the vertex correction). The vacuum-polarization modification of the photon propagator (Section 3.1) is one of three one-loop QED corrections. The vertex correction — a virtual photon emitted and reabsorbed by the electron line — modifies the electron’s magnetic moment to $a_e = (g-2)/2 = \alpha/(2\pi) \approx 0.00116$ (Schwinger, 1948). This is a dimension-5 effective operator $\bar{\psi} \sigma^{\mu\nu} \psi F_{\mu\nu}$ in the EFT language of Section 4 — the magnetic analog of the electric contact expansion. Higher-order corrections involve both vacuum polarization insertions (the running coupling of Section 3.1) and light-by-light fermion loops, and the QED prediction agrees with the measured electron $g-2$ to better than one part in 10^{12} — the most precisely tested prediction in physics. The muon anomalous magnetic moment, being more sensitive to heavy virtual particles by a factor $\sim (m_\mu/m_e)^2$, probes the contact operators of heavier sectors: new physics at scale Λ contributes $\delta a_\mu \sim m_\mu^2/\Lambda^2$, the standard EFT suppression pattern.

Remark 3.4 (Schwinger pair production and the Euler–Heisenberg Lagrangian: the non-perturbative complement). The perturbative vacuum-polarization story of Section 3.1 has a non-perturbative counterpart. In a constant electric field \mathcal{E} , the QED vacuum is unstable against spontaneous electron–positron pair creation at a rate per unit volume $\Gamma/V \propto m^2(\mathcal{E}/\mathcal{E}_{\text{cr}})^2 \exp(-\pi\mathcal{E}_{\text{cr}}/\mathcal{E})$, where $\mathcal{E}_{\text{cr}} = m^2c^3/(e\hbar) \approx 1.3 \times 10^{18} \text{ V/m}$ is the Schwinger critical field (Schwinger, 1951). The essential singularity $\exp(-\pi\mathcal{E}_{\text{cr}}/\mathcal{E})$ makes this invisible to all orders of perturbation theory: the rate vanishes faster than any power of α as $\mathcal{E} \rightarrow 0$. The real part of the same one-loop effective action gives the Euler–Heisenberg Lagrangian $\mathcal{L}_{\text{EH}} = \mathcal{L}_{\text{Maxwell}} + (2\alpha^2/45m^4)[(\mathcal{E}^2 - \mathcal{B}^2)^2 + 7(\mathcal{E} \cdot \mathcal{B})^2] + \dots$ (Euler and Heisenberg, 1936) — dimension-8 contact operators in the electromagnetic field, the QFT counterpart of the contact expansion in Section 4. Fermions thus affect forces both perturbatively (running coupling, Uehling potential) and non-perturbatively (pair creation, photon–photon scattering), with the contact-operator tower organizing both regimes.

Remark 3.5 (Dispersion relations: the analytic bridge between running coupling and pair creation). The vacuum polarization function $\Pi(q^2)$ of Section 3.1 is analytic in q^2 except for a branch cut starting at $q^2 = 4m_e^2$ (the pair-creation threshold of Remark 3.4). The Kramers–Kronig dispersion relation (Kramers, 1927; Kronig, 1926) — expressing causality as analyticity in the upper half-plane — relates the real and imaginary parts: $\Pi(q^2) - \Pi(0) = (q^2/\pi) \int_{4m_e^2}^\infty ds \text{Im } \Pi(s)/[s(s - q^2)]$, where the subtraction at $q^2 = 0$ absorbs the charge renormalization. The imaginary part $\text{Im } \Pi(s)$ is the absorptive cross-section for the corresponding process (e.g., $\sigma(e^+e^- \rightarrow \text{hadrons})$ for hadronic vacuum polarization, via the optical theorem), so the real part — the running coupling of Section 3.1 — is entirely determined by production thresholds and cross-sections. This is the “contact data” theme of Section 4 in analytic dress: the function is determined by its singularities (the branch cut), and the dispersion integral is a controlled reconstruction from that singular support, paralleling the way the effective range expansion of Remark 4.2 encodes scattering observables in a few parameters.

Remark 3.6 (Optical theorem: unitarity as the physical engine behind dispersion). The dispersion reconstruction of Remark 3.5 rests on analyticity (causality), but the physical content of the spectral weight $\text{Im } \Pi(s)$ comes from unitarity. The optical theorem — a direct consequence of $SS^\dagger = 1$ — states that the imaginary part of the forward elastic scattering amplitude is proportional to the total cross-section (summing over all open channels, elastic and inelastic). For the vacuum polarization, this connects $\text{Im } \Pi(s)$ to the cross-section $\sigma(e^+e^- \rightarrow X)$ at center-of-mass energy \sqrt{s} , through the specific structure of the photon self-energy: the photon propagator dressed by the fermion loop mediates the process, and its imaginary part above the pair-production threshold $s = 4m^2$ is the production rate. At the diagrammatic level, Cutkosky’s cutting rules (1960) compute the discontinuity by replacing cut propagators with on-shell delta functions $\delta(p^2 - m^2)\theta(p^0)$, putting intermediate

states on-shell — each cut corresponds to a physical production channel in the unitarity sum. The structural point for Section 4’s contact expansion is that the polynomial $C_0 + C_2q^2 + \dots$ is manifestly real below threshold; the non-polynomial, imaginary content that appears above threshold (open production channels) is precisely what the contact expansion cannot capture — it is the absorptive physics that the dispersion integral reconstructs from the branch-cut singularity.

Remark 3.7 (Schwinger model: exact fermion-to-boson transmutation in 1+1 dimensions). The Schwinger model — QED in 1+1 dimensions with massless fermions coupled to a U(1) gauge field (Schwinger, 1962) — is exactly solvable via bosonization (Coleman, 1975) and provides the simplest demonstration of how fermionic fields produce a purely bosonic physical spectrum [Schwinger1962Mass]. The model confines: the photon acquires a mass $m = e/\sqrt{\pi}$ (where e is the gauge coupling, with mass dimension 1 in 1+1D), and the physical spectrum consists of a single massive boson — the “Schwinger boson,” a fermion-antifermion composite. No free fermions appear as asymptotic states: the exact fermion propagator has its pole completely cancelled by gauge interactions [ColemanJackiwSusskind1975]. The chiral anomaly plays the central role: the axial current satisfies $\partial_\mu j^{\mu 5} = (e/(2\pi)) \varepsilon^{\mu\nu} F_{\mu\nu}$, and it is this anomaly that generates the boson mass — the mass formula $m^2 = e^2/\pi$ is exact to all orders, receiving no higher-loop corrections. For the massive Schwinger model (fermion mass $m_f \neq 0$), a non-trivial θ -vacuum structure appears: the topological vacuum angle produces a background electric field, giving the 1+1D analog of the QCD vacuum structure of Remark 6.4. The model demonstrates the full fermion \rightarrow boson chain of Section 6 in an exactly solvable setting: fermions are permanently confined into bosonic composites, the mass scale arises entirely from the anomaly (not from an explicit fermion mass term in the Lagrangian), and at distances $r \gg 1/m$ the interaction is exponentially screened, collapsing to the contact expansion of Section 4.

4. IR replacement: local operators \Rightarrow contact kernels / boundary-condition data

At low resolution, integrating out heavy degrees of freedom typically produces local operators. In a two-body, nonrelativistic sector, this appears as an amplitude expansion analytic in momentum transfer q :

$$\mathcal{A}(q) = C_0 + C_2 q^2 + O(q^4).$$

The coordinate-space interaction associated to such an analytic expansion is distributional and diagonal-supported. The invariant core is a Fourier-transform identity:

$$\int \frac{d^d q}{(2\pi)^d} e^{iq \cdot r} = \delta^{(d)}(r), \quad \int \frac{d^d q}{(2\pi)^d} q^2 e^{iq \cdot r} = -\nabla^2 \delta^{(d)}(r).$$

Derivation 4.1 (Contact expansion gives \(\delta\)-functions and derivative contacts). Interpreting the low-energy interaction kernel as the inverse Fourier transform of $\mathcal{A}(q)$ (Born-level language, up to overall convention-dependent factors), we obtain

$$V(r) \propto \int \frac{d^d q}{(2\pi)^d} e^{iq \cdot r} \mathcal{A}(q) \propto C_0 \delta^{(d)}(r) - C_2 \nabla^2 \delta^{(d)}(r) + \dots$$

Thus locality at low energy naturally becomes **contact data**: delta kernels and their derivatives supported at a point (or on the diagonal, in kernel language).

In singular cases (notably δ interactions in $d \geq 2$ in various channels), this contact data is not simply a fixed number: it is defined by a renormalization condition and can generate RG-invariant scales and bound states [Jackiw1991DeltaPotentials] [ManuelTarrach1994PertRenQM].

Remark 4.2 (Connection to the effective range expansion). In scattering theory the s-wave amplitude is parametrized by the effective range expansion (ERE) $k \cot \delta_0(k) = -1/a + r_0 k^2/2 + O(k^4)$, where a is the scattering length and r_0 the effective range. The contact expansion of Derivation 4.1 is the momentum-space counterpart: at Born level, C_0 determines a , C_2 determines r_0 , and each higher C_{2n} maps to a shape parameter. Examples 5.1 and 5.2 below are the leading-order case $C_2 = 0$ (zero effective range, $r_0 = 0$), for which the full amplitude $f_0(k) = -a/(1 + ika)$ depends on a single parameter — the scattering length.

Remark 4.3 (Huang-Yang pseudopotential: regularization built into the operator). The Huang-Yang pseudopotential $V(r) = (4\pi\hbar^2 a/m) \delta^{(3)}(r) (\partial/\partial r)(r \cdot)$ builds the renormalization condition of Example 5.2 directly into the operator definition. The differential operator $(\partial/\partial r)(r \cdot)$ extracts the regular part of the wavefunction at the origin — for a function with the s-wave asymptotics $\psi(r) \sim A(1/r - 1/a) + (\text{regular})$, it yields the finite value $-A/a$ rather than the divergent A/r — and thereby automatically implements the scattering-length boundary condition without explicit cutoff manipulation. This is a “smart” contact kernel: the prescription for handling the $r = 0$ singularity is not an external regularization step but is part of the definition of the interaction, with different values of a selecting different self-adjoint extensions.

Remark 4.4 (Weinberg power counting: the systematic hierarchy for contact EFTs). The contact expansion of Derivation 4.1 is the leading term in a systematic effective field theory organized by Weinberg’s power counting (1990, 1991). Each diagram is assigned a chiral order based on the number of loops L , the number of derivatives d_i at each vertex, and the number of nucleon lines n_i — with higher-order contributions suppressed by powers of Q/Λ , the ratio of typical momentum to the breakdown scale. At leading order (LO), only the non-derivative contact C_0 contributes; at next-to-leading order (NLO), the two-derivative contact $C_2 q^2$ and one-loop diagrams with C_0 insertions enter together; at NNLO, $C_4 q^4$, mixed one-loop diagrams, and two-loop diagrams

appear. Crucially, loops generate non-analytic momentum dependence — terms like $q^2 \ln(q^2/\mu^2)$ in three dimensions — that cannot be absorbed into the polynomial contact expansion. These infrared logarithms distinguish the full quantum EFT from a naive polynomial fit and provide the physical content (e.g., the running coupling of Examples 5.1–5.2) that transcends tree-level contact operators. When the leading contact interaction is strong (large scattering length), an alternative counting due to Kaplan, Savage, and Wise (1998) resums C_0 to all orders — precisely the T -matrix resummation of Examples 5.1–5.2 — treating higher-order contacts and pion exchanges as perturbative corrections.

Remark 4.5 (Operator product expansion: the QFT generalization of the contact hierarchy). Wilson’s operator product expansion (OPE, 1969) extends the contact hierarchy of Derivation 4.1 to relativistic quantum field theory. When two local operators $\mathcal{O}_A(x)$ and $\mathcal{O}_B(0)$ are brought together inside time-ordered correlation functions, their product admits an expansion in local operators at the coincidence point: $\mathcal{O}_A(x)\mathcal{O}_B(0) \sim \sum_k C_k^{AB}(x)\mathcal{O}_k(0)$, where the Wilson coefficient functions scale as $C_k \sim |x|^{d_k - d_A - d_B}$ (with d_i the engineering dimensions), up to logarithmic corrections from anomalous dimensions in non-conformal theories. The identity-operator contribution ($\mathcal{O}_k = \mathbf{1}$) yields a c-number equal to the interacting two-point function (reducing to the propagator K^{-1} of Section 2 in the free-field limit), while subleading operators are suppressed by positive powers of $|x|$ times a physical mass scale — the same hierarchy principle as the $C_0 + C_2 q^2 + \dots$ contact expansion, though the two operate in opposite regimes: the OPE is a short-distance (UV) expansion ($|x| \rightarrow 0$), while the contact expansion is a low-momentum (IR) expansion ($q \rightarrow 0$). The structural parallel lies in the shared organizational principle: unresolved physics is encoded by a tower of local operators ordered by dimension. The OPE is a general structural feature of local QFT — in conformal field theories (no mass gap) it is exact and convergent, constituting the fundamental algebraic data of the theory; in massive theories it is an asymptotic short-distance expansion. The scale dependence of OPE coefficients in non-conformal theories is governed by the Callan-Symanzik equation, whose beta function is the direct QFT counterpart of the contact-coupling beta function in Examples 5.1–5.2.

5. Boundary-condition reading (point interactions)

Point-supported interactions can be encoded as self-adjoint extension / boundary-condition data rather than as ordinary functions $V(r)$. This is the natural operator-theoretic counterpart of “diagonal-support kernels.” For standard references and pedagogical framing, see [BonneauFarautValent2001SAE] and the delta-potential discussion in [Jackiw1991DeltaPotentials].

This viewpoint matches the controlled-refinement perspective: when a continuum

description is defined as a refinement limit, UV data can survive in the limit precisely as boundary-condition parameters (contact terms), with RG flow expressing compatibility across resolutions.

Example 5.1 (Contact coupling generates a scale: 2D delta potential). In two spatial dimensions, a contact interaction $V(r) = g_0 \delta^{(2)}(r)$ with bare coupling g_0 and UV cutoff Λ leads, after a standard loop integral, to a T -matrix with the structure

$$T(k)^{-1} = \frac{1}{g_0} + \frac{m}{\pi \hbar^2} \ln\left(\frac{\Lambda}{k}\right),$$

which diverges as $\Lambda \rightarrow \infty$ unless g_0 is tuned. Define a renormalized coupling at reference scale μ by absorbing the $\ln \Lambda$ divergence; cutoff independence then gives the beta function $\beta(g_R) = \mu dg_R/d\mu = (m/\pi \hbar^2) g_R^2$. This is a quadratic beta function of the same form as the toy logarithmic model in the cornerstone (Section 8.3), with solution producing a dynamically generated scale $\mu_* = \mu e^{\pi \hbar^2/(mg_R)}$. For attractive coupling ($g_R < 0$) this scale is below the reference scale and sets the bound-state energy: $E = -\hbar^2 \mu_*^2/(2m)$.

The structural lesson: the “contact” limit of the effective interaction is not a number (coupling constant) but a flow — a scale-dependent parameter whose RG trajectory is part of the definition. This is “uncuttable” in the sense of the companion note: the continuum theory requires the refinement rule (cutoff removal + beta function) and not merely a single-cutoff value.

Example 5.2 (3D contact interaction: scattering length). In three spatial dimensions, the same contact interaction $V(r) = g_0 \delta^{(3)}(r)$ with UV cutoff Λ produces a linearly divergent loop integral (compared to the logarithmic divergence in $d = 2$). After resummation, the s-wave scattering amplitude takes the standard effective-range form with zero effective range:

$$f_0(k) = \frac{-a}{1 + ika},$$

where the scattering length a is defined by absorbing the Λ -dependent bare coupling into a single physical parameter via a renormalization condition of the form $1/g_0 \propto \Lambda + (\text{finite part depending on } a)$. When $a > 0$, a pole at $k = i/a$ gives a bound state with energy $E = -\hbar^2/(2ma^2)$ [AlbeverioGesztesyHoeghKrohnHolden2005]. The comparison with Example 5.1 highlights how the divergence character changes with dimension — logarithmic ($d = 2$) versus linear ($d = 3$) — while the structural lesson is identical: the “coupling constant” of a contact interaction is not a bare number but a renormalization-group datum, defined only through a refinement rule (cutoff removal + physical matching condition).

Remark 5.3 (Unitarity limit: universality at the RG fixed point). When the scattering length diverges ($|a| \rightarrow \infty$), the contact coupling sits at a non-trivial RG fixed point and the theory becomes scale-invariant: no microscopic length survives beyond the interparticle spacing. Thermodynamic ratios become universal — for a spin- $\frac{1}{2}$ Fermi gas the ground-state energy is

$E = \xi E_{\text{FG}}$ with Bertsch parameter $\xi \approx 0.37$, independent of the short-range physics that produced the large scattering length. This fixed point controls the BEC–BCS crossover in cold atomic gases, where a magnetic Feshbach resonance tunes a through $\pm\infty$, providing a laboratory realization of the structural lesson in Examples 5.1–5.2: the contact coupling is not a number but a flow, and the fixed point of that flow generates universality.

Remark 5.4 (Dimensional dependence: from UV-finite to cutoff-dependent extension data). The same contact operator $\delta^{(d)}(r)$ requires qualitatively different control data across dimensions. In $d = 1$, the delta potential defines a self-adjoint extension of the free Hamiltonian whose boundary-condition parameter (the coupling g) is UV-finite — no cutoff dependence, no RG flow, just a number. In $d = 2$ (Example 5.1), the extension parameter diverges logarithmically with the cutoff, producing a running coupling and dimensional transmutation. In $d = 3$ (Example 5.2), the divergence is linear. The critical dimension is $d = 2$, where the contact coupling is classically marginal; above it, the coupling is classically relevant and power-law subtractions are needed. In all dimensions the interaction requires self-adjoint extension theory (it is never a “plain operator perturbation”), but the extension datum transitions from UV-finite ($d = 1$) to cutoff-dependent ($d \geq 2$) — and it is this transition that makes renormalization group flow part of the definition.

Remark 5.5 (Efimov effect: 3-body limit cycle from 2-body contact). At the unitarity limit of Remark 5.3, the 3-body sector exhibits a qualitatively new RG phenomenon: the running 3-body coupling is log-periodic in the cutoff with period π/s_0 (where $s_0 \approx 1.006$ for identical bosons), corresponding to a geometric energy scaling factor $e^{\pi/s_0} \approx 22.7$. This RG limit cycle — as opposed to the 2-body fixed point — produces an infinite tower of 3-body bound states with geometrically spaced energies $E_n \propto e^{-2\pi n/s_0}$, a signature of discrete scale invariance first predicted by Efimov (1970) and observed in cesium cold atoms by Kraemer et al. (2006). Crucially, the absolute position of the tower requires an additional 3-body parameter beyond the scattering length a — a new piece of self-adjoint extension data not determined by 2-body physics, illustrating that the contact/extension hierarchy grows richer with particle number.

Remark 5.6 (Feshbach resonances: laboratory realization of the contact-coupling RG flow). The theoretical framework of Examples 5.1–5.2 has a direct experimental counterpart: a magnetic Feshbach resonance occurs when a bound state in a closed hyperfine channel is Zeeman-tuned into near-degeneracy with the open scattering threshold, producing a scattering length $a(B) = a_{\text{bg}}(1 - \Delta B/(B - B_0))$ that diverges as $B \rightarrow B_0$ (Chin et al., 2010). Sweeping B through the resonance traces the full one-parameter family of self-adjoint extensions: from weakly attractive ($a < 0$, BCS pairing) through the scale-invariant unitarity fixed point of Remark 5.3 ($|a| = \infty$) to the BEC regime of tightly bound dimers ($a > 0$), thereby realizing the BEC–BCS crossover as a continuous traversal of the RG flow. The Efimov states of Remark 5.5 have been observed near such Feshbach resonances in cesium (Kraemer et al.,

2006). Structurally, the magnetic field B is an external knob that tunes the renormalized contact coupling at a fixed energy scale — the experimentalist’s version of the beta function controlling the coupling flow in Examples 5.1–5.2.

Remark 5.7 (Casimir effect: macroscopic force from boundary conditions alone). The self-adjoint-extension/boundary-condition paradigm of Examples 5.1–5.2 generates a measurable macroscopic force in the Casimir effect (Casimir, 1948). Two perfectly conducting parallel plates separated by distance d impose boundary conditions that restrict the electromagnetic vacuum modes to a discrete set between the plates. The regulated difference of zero-point energies — computed via zeta-function regularization, exponential cutoff, or dimensional methods, all yielding the same result — gives an attractive force per unit area $F/A = -\pi^2 \hbar c/(240 d^4)$, experimentally verified to percent-level precision (Lamoreaux, 1997; Mohideen and Roy, 1998). No classical source $J(x)$ is needed: the force arises purely from how the boundary conditions modify the vacuum fluctuation spectrum. The regularization-scheme independence of the result echoes the cutoff independence of the renormalized observables in Examples 5.1–5.2: the physically meaningful Casimir energy is the refinement limit (regulator removed), not any single regulated approximant. This is the boundary-condition paradigm of Section 5 elevated from a mathematical framework to a laboratory-observable phenomenon.

Remark 5.8 (Levinson’s theorem: boundary-condition data counts bound states). The self-adjoint-extension framework of Examples 5.1–5.2 organizes both scattering and bound-state data through a single boundary-condition parameter a . Levinson’s theorem (1949) makes this spectral connection precise: for an s-wave ($l = 0$) potential with sufficient decay at infinity (short-range), the phase shift satisfies $\delta_0(0) = n_0\pi$, where n_0 is the number of s-wave bound states and the convention $\delta_0(k \rightarrow \infty) = 0$ is adopted. The total phase accumulated from high to low energy exactly counts the bound states — each bound state contributes π of phase. When a zero-energy resonance (half-bound state) exists — a solution at $k = 0$ that is finite but not square-integrable — the relation becomes $\delta_0(0) = (n_0 + \frac{1}{2})\pi$. For the contact interaction of Example 5.2, this half-bound-state condition corresponds precisely to $|a| = \infty$, the unitarity limit of Remark 5.3: the divergent scattering length signals the threshold between having and not having a bound state. The theorem thus connects the infrared boundary-condition datum a (the effective range expansion of Remark 4.2) to the spectral content (bound-state count) of the self-adjoint extension.

Remark 5.9 (Tan’s contact: many-body thermodynamics from boundary-condition data). The boundary-condition parameter a of Examples 5.1–5.2 governs not only two-body scattering but also universal properties of many-body systems. Tan (2008) showed that for a two-component Fermi gas with s-wave contact interactions, the short-range pair structure is characterized by a single extensive quantity \mathcal{C} (the “contact”) through a family of exact relations: the momentum distribution obeys $n_\sigma(k) \rightarrow \mathcal{C}/k^4$ as $k \rightarrow \infty$ (the k^{-4} tail being the

momentum-space signature of the contact interaction of Section 4); the adiabatic sweep theorem $dE/d(-1/a) = \hbar^2\mathcal{C}/(4\pi m)$ relates the energy response to scattering-length variation; and for a homogeneous system the pressure satisfies $P = \frac{2}{3}(E/V) + \hbar^2\mathcal{C}/(12\pi maV)$. The contact \mathcal{C} is a many-body observable that is sensitive to the self-adjoint-extension parameter a : it measures the number of pairs at short range (the opposite-spin pair density diverges as $\rho_2(r \rightarrow 0) \sim \mathcal{C}/(4\pi r)^2$), thereby elevating the two-body boundary-condition datum to a thermodynamic quantity. These relations have been verified experimentally in ultracold Fermi gases (Stewart et al., 2010), where the momentum-distribution tail and the equation of state were measured in the same system across the BEC–BCS crossover of Remark 5.6.

6. Outlook: emergent bosonic composites

Remark 6.1 (Cooper pairing: fermion bilinears as emergent bosonic mediators). The Grassmann-source obstruction of Section 3 is bypassed when fermion bilinears condense into bosonic composite modes. The canonical example is BCS superconductivity (Bardeen, Cooper, and Schrieffer, 1957): any attractive interaction in the Cooper channel between fermions near a Fermi surface produces a pairing instability with a non-perturbative gap $\Delta \propto \omega_D \exp(-1/(N(0)|V|))$, where $N(0)$ is the density of states and $|V|$ the interaction strength. The resulting condensate — described by the bosonic order parameter $\langle \psi_\downarrow \psi_\uparrow \rangle$ — gives mass to the photon via the Anderson–Higgs mechanism, with the effective photon mass $m_A \propto \sqrt{n_s e^2}$ set by the superfluid density n_s (proportional to $|\Delta|^2$ at zero temperature). This generates the Meissner effect: magnetic fields decay exponentially inside the superconductor with penetration depth $\lambda = 1/m_A$. In the contact-interaction limit, the BCS gap equation requires the same renormalization as Examples 5.1–5.2: the bare coupling diverges with the cutoff, and the physical gap is determined by the renormalized scattering length. The BEC–BCS crossover of Remark 5.6 then interpolates continuously between overlapping Cooper pairs (BCS) and tightly bound bosonic molecules (BEC), demonstrating that the “emergent boson” spans the full range from collective many-body mode to genuine two-body bound state.

Remark 6.2 (Neutrino-mediated forces: the direct fermionic potential). The paper’s motivating question — “can a fermionic field generate a potential?” — has a direct quantitative answer. Feinberg and Sucher (1968) computed the long-range force between fermions mediated by virtual neutrino-antineutrino pair exchange (a box diagram with two neutrino propagators and two weak vertices). The result is $V(r) \propto G_F^2/r^5$ — a power-law potential steeper than the Coulomb $1/r$ or Yukawa e^{-mr}/r of bosonic exchange, with the $1/r^5$ scaling reflecting the two-fermion intermediate state. Crucially, this is not a tree-level exchange (which would require Grassmann sources, per Remark 3.1) but a loop effect: two fermion propagators carrying the force, exactly as Section

3.1's controlled alternative predicts. The power law (rather than exponential suppression) arises because the neutrino is (nearly) massless; for massive neutrinos, the potential crosses over to exponential decay at distances $r \gg 1/m_\nu$. At short distances $r \ll 1/M_W$, the two W -boson propagators collapse to the Fermi four-fermion vertex, and the force reduces to the contact expansion of Section 4 — the neutrino-mediated potential is the long-range tail of a contact interaction viewed at low resolution.

Remark 6.3 (Nuclear forces from pion exchange: the paradigmatic fermion-to-boson-to-potential chain). The Grassmann-source obstruction of Section 3 is bypassed most famously in the nuclear force. Quarks (fermions) bind into pions — pseudoscalar bosonic composites with the quantum numbers of $\bar{q}\gamma_5\tau q$ (where τ are isospin Pauli matrices) — which are the pseudo-Goldstone bosons of spontaneously broken chiral symmetry. Once formed, pions mediate the longest-range part of the nucleon-nucleon interaction through one-pion exchange (OPEP), giving a Yukawa potential $V(r) \propto e^{-m_\pi r}/r$ with range $\sim 1/m_\pi \approx 1.4$ fm, as predicted by Yukawa (1935). The OPEP includes a tensor component (proportional to the spin-dependent operator $S_{12} = 3(\boldsymbol{\sigma}_1 \cdot \hat{r})(\boldsymbol{\sigma}_2 \cdot \hat{r}) - \boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2$) responsible for the deuteron quadrupole moment and D -state admixture. At distances shorter than the pion Compton wavelength, the pion-exchange potential is not resolved and collapses into the contact expansion of Section 4 — the same structural pattern as the neutrino-mediated force of Remark 6.2. Modern chiral effective field theory (Weinberg, 1990; Epelbaum, Hammer, and Meißner, 2009) organizes the nuclear force as a systematic hierarchy: pion exchanges provide the long-range, non-analytic-in-momentum contributions, while short-range physics is encoded by the contact operators C_0, C_2, \dots of Derivation 4.1, with the expansion parameter Q/Λ_χ (where $\Lambda_\chi \sim 1$ GeV is the chiral symmetry breaking scale). This is the complete realization of the paper's structural chain: fermions (quarks) \rightarrow emergent bosons (pions) \rightarrow Yukawa potential (long range) + contact operators (short range), with the EFT framework providing the systematic power counting that relates the two regimes.

Remark 6.4 (Color confinement: the dynamical completion of the fermion-to-composite chain). The structural chain of Section 6 — fermions \rightarrow emergent composites \rightarrow effective interactions — reaches its most dramatic realization in QCD. Quarks interact via gluon exchange, but at long range the quark-antiquark potential grows linearly: $V(r) \sim \sigma r$, with string tension $\sigma \approx 0.18$ GeV 2 (in natural units, equivalently ~ 0.9 GeV/fm). Wilson's area-law criterion (1974) makes this precise: for a rectangular $R \times T$ Wilson loop, $\langle W(C) \rangle \sim \exp(-\sigma RT)$ signals confinement (linear potential), versus $\langle W(C) \rangle \sim \exp(-\mu P)$ (perimeter law, $P = 2(R + T)$) for a deconfined/screened phase. When the flux-tube energy becomes sufficient to materialize a new quark-antiquark pair from the vacuum, the string fragments — producing hadrons rather than free quarks. Confinement thus produces both bosonic color-singlet composites (mesons $q\bar{q}$, glueballs) and fermionic ones (baryons qqq); the structural point for this paper is that all long-range hadronic interactions are

mediated by the bosonic composites (pion exchange of Remark 6.3), completing the chain from fundamental fermions to effective bosonic potentials. The pattern echoes — without being logically equivalent to — the Grassmann-source obstruction of Section 3: the fundamental fermionic degrees of freedom are not directly accessible as asymptotic states, and the effective low-energy description (chiral perturbation theory) is built entirely from bosonic meson fields. The large- N expansion ('t Hooft, 1974), in which planar gluon diagrams dominate, provides complementary analytical evidence for string-like behavior of the confining flux tube. Proving that pure Yang-Mills gauge theory in four dimensions has a mass gap from first principles is one of the Clay Millennium Prize Problems; lattice QCD provides numerical evidence for confinement and a mass gap, but a rigorous proof remains open.

Remark 6.5 (Hubbard-Stratonovich transformation: making the emergent boson manifest). The emergent-boson picture of Remarks 6.1–6.3 can be made systematic at the path-integral level via the Hubbard-Stratonovich (HS) transformation (Hubbard, 1959; Stratonovich, 1957). Starting from an attractive four-fermion interaction $S_{\text{int}} = -g \int (\bar{\psi}\psi)^2$ (with $g > 0$), a Gaussian identity introduces an auxiliary scalar field σ that linearizes the quartic coupling: $\exp(g \int (\bar{\psi}\psi)^2) = \int \mathcal{D}\sigma \exp(-\int [\sigma^2/(4g) - \sigma\bar{\psi}\psi])$. The quartic fermion interaction becomes a Yukawa coupling $\sigma\bar{\psi}\psi$, and the fermion path integral — now Gaussian in $\psi, \bar{\psi}$ — evaluates to a functional determinant, giving the bosonic effective action $S_{\text{eff}}[\sigma] = \int \sigma^2/(4g) - \text{Tr} \ln(\partial - m - \sigma)$. The trace-log encodes all fermion-loop contributions to the σ dynamics; its saddle-point approximation (controlled when the number of fermion species N_f is large) gives mean-field results. In the BCS channel (Remark 6.1), the HS field is the gap parameter Δ , and the saddle-point condition reproduces the BCS gap equation; fluctuations around the saddle yield the Anderson–Bogoliubov sound mode and the Higgs amplitude mode. In the Nambu–Jona-Lasinio model (Nambu and Jona-Lasinio, 1961), the same construction introduces σ and π fields from quartic quark interactions; the saddle point $\langle\sigma\rangle \neq 0$ realizes chiral symmetry breaking, with pions as the pseudo-Goldstone bosons of Remark 6.3. A key caveat: the HS transformation requires choosing a specific interaction channel (particle-particle for BCS, particle-hole for chiral condensation), via Fierz rearrangement of the four-fermion vertex; different channel choices produce different auxiliary fields and different mean-field physics. The RG flow of the original fermionic theory is encoded in the σ effective action through the trace-log: integrating out high-momentum fermion shells generates running couplings for σ , with matching at the compositeness scale connecting the fermionic UV description to the bosonic IR effective theory — the same structural chain as the pion-exchange story of Remark 6.3.

Remark 6.6 ('t Hooft anomaly matching: the topological constraint on the fermion-to-composite chain). The fermion → emergent boson chain of Section 6 is constrained by anomaly matching ('t Hooft, 1980): if a global symmetry G has an anomaly in the UV theory (computed from the fundamental fermion content), then the same anomaly must appear in the IR

theory — either from massless composite fermions, from Goldstone bosons (via the Wess-Zumino-Witten term), or from symmetry-protected topological phases [’tHooft1980]. The anomaly coefficients are discrete (related to topological invariants via the Atiyah-Singer index theorem) and cannot vary under continuous RG flow, making ’t Hooft anomalies for global symmetries exact RG invariants [WessZumino1971]. For QCD (Remarks 6.3–6.4), the fundamental theory has N_f quarks giving anomaly coefficients for $SU(N_f)_L \times SU(N_f)_R \times U(1)_B$; with chiral symmetry breaking $SU(N_f)_L \times SU(N_f)_R \rightarrow SU(N_f)_V$, the pions (Goldstone bosons) carry the anomaly through the Wess-Zumino-Witten topological term (Witten, 1983) [Witten1983WZW]. The quantitative test: the $\pi^0 \rightarrow \gamma\gamma$ decay rate, completely fixed by the anomaly coefficient (proportional to N_c , the number of colors), agrees with experiment — directly verifying that pions carry the quark anomaly. When the Hubbard-Stratonovich auxiliary field of Remark 6.5 is interpreted as a dynamical IR degree of freedom, anomaly matching constrains which channels can consistently realize the fundamental symmetries: not every Fierz channel produces an anomaly-compatible low-energy theory. Anomaly matching thus provides a non-perturbative selection rule on the fermion-to-composite chain: the composites are not arbitrary but must be “anomaly-compatible” with the fundamental fermions.

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