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# The Lamb Shift from an Effective Hamiltonian in Non-relativistic Quantum Electrodynamics and a General Class of Effective Operators

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## Abstract

Some aspects of spectral analysis of an effective Hamiltonian in non-relativistic quantum electrodynamics are reviewed. The Lamb shift of a hydrogen-like atom is derived as the lowest order approximation (in the fine structure constant) of an energy level shift of the effective Hamiltonian. Moreover, a general class of effective operators is presented, which comes from models of an abstract quantum system interacting with a Bose field.

**Keywords:** Non-relativistic quantum electrodynamics, effective Hamiltonian, spectrum, Lamb shift, effective operator

**Mathematics Subject Classification 2010:** 47N50, 81Q10, 81Q15, 81V10

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## 1 Introduction

This paper is a review of some results obtained in [7], in which spectral analysis is made on an effective Hamiltonian in non-relativistic quantum electrodynamics (QED), a quantum theory of non-relativistic charged particles interacting with the quantum radiation field (a quantum field theoretical version of a vector potential in classical electrodynamics). We also present a general class of effective operators for an abstract quantum system interacting with a Bose field (a quantum field of bosons) such that the effective Hamiltonian in non-relativistic QED is a concrete example in the class.

In this introduction, we first explain some physical backgrounds behind the work [7].

A hydrogen-like atom is an atom consisting of one electron, whose electric charge is  $-e < 0$ , and a nucleus with electric charge  $Ze > 0$ , where  $Z$  is a natural number (the case  $Z = 1$  is the usual hydrogen atom). As is well known, if the nucleus is fixed at the origin of the 3-dimensional Euclidean vector space  $\mathbb{R}^3 = \{\mathbf{x} = (x_1, x_2, x_3) | x_j \in \mathbb{R}, j = 1, 2, 3\}$  and, as the potential acting on the electron at the position  $\mathbf{x} \in \mathbb{R}^3$ , one takes into account only the electric Coulomb potential<sup>1</sup>  $-Ze^2/4\pi|\mathbf{x}|$  from the nucleus, then a quantum mechanical Hamiltonian describing the hydrogen-like atom is given by the Schrödinger operator

$$H_{\text{hyd}} = -\frac{\hbar^2}{2m_e}\Delta - \frac{\gamma}{|\mathbf{x}|} \quad (1.1)$$

acting on  $L^2(\mathbb{R}^3)$ , the Hilbert space of equivalence classes of complex-valued functions square integrable on  $\mathbb{R}^3$  with respect to the 3-dimensional Lebesgue measure, where  $\hbar := h/2\pi$  ( $h$  is the Planck constant),  $m_e > 0$  is the electron mass,  $\Delta$  is the generalized Laplacian on  $L^2(\mathbb{R}^3)$ , and

$$\gamma := \frac{Ze^2}{4\pi}.$$

Indeed,  $H_{\text{hyd}}$  is self-adjoint with domain  $D(H_{\text{hyd}}) = D(\Delta)$ —for a linear operator  $A$  on a Hilbert space,  $D(A)$  denotes the domain of  $A$ —and the spectrum of  $H_{\text{hyd}}$ , denoted  $\sigma(H_{\text{hyd}})$ , is found to be

$$\sigma(H_{\text{hyd}}) = \{E_n\}_{n=1}^{\infty} \cup [0, \infty) \quad (1.2)$$

with

$$E_n = -\frac{1}{2} \frac{m_e \gamma^2}{\hbar^2} \frac{1}{n^2}, \quad n = 1, 2, 3, \dots, \quad (1.3)$$

where each eigenvalue  $E_n$  is degenerate with multiplicity  $n^2$  (e.g., [6, §2.3.5a] and [6, Lemma 5.22, footnote 12]). These eigenvalues explain very well the so-called principal

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<sup>1</sup>The electromagnetic system of units which we use in the present paper is the rationalized CGS Gauss unit system with the dielectric constant in the vacuum equal to 1.

energy levels of the hydrogen-like atom (Fig.1(a)), but do not show the finer structures of the energy spectrum (Fig.1(b)), which may be regarded as splits of the degeneracy of  $E_n$ 's .

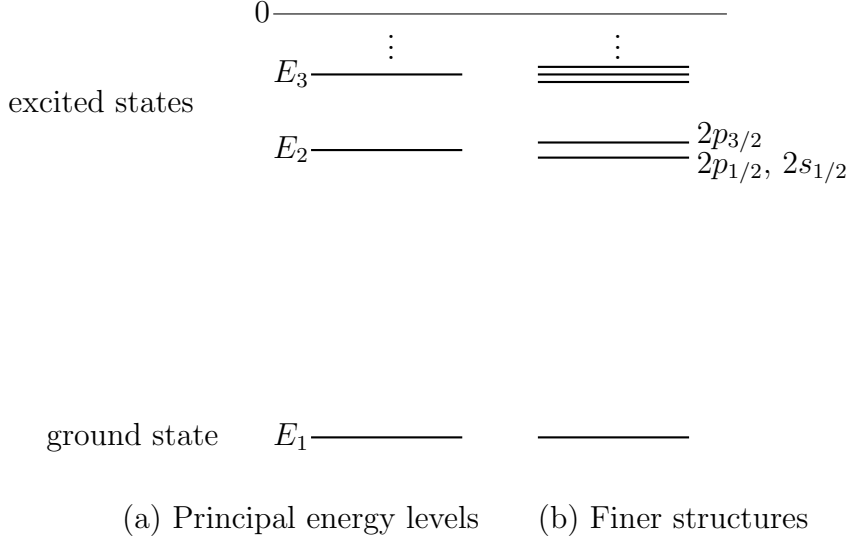


Figure 1: Energy spectrum of a hydrogen-like atom

It turns out that the finer structures of the hydrogen-like atom can be explained by the Dirac operator

$$D_{\text{hyd}} := -i\hbar c \sum_{k=1}^3 \alpha_k D_k + m_e c^2 \beta - \frac{\gamma}{|\mathbf{x}|},$$

acting on  $\oplus^4 L^2(\mathbb{R}^3)$  (the four direct sum of  $L^2(\mathbb{R}^3)$ ), where  $c > 0$  is the speed of light in the vacuum,  $D_k$  is the generalized partial differential operator in the variable  $x_k$ , and  $\alpha_k, \beta$  are  $4 \times 4$  Hermitian matrices satisfying the following anti-commutation relations ( $\delta_{kl}$  denotes the Kronecker delta):

$$\alpha_k \alpha_l + \alpha_l \alpha_k = 2\delta_{kl}, \quad \alpha_k \beta + \beta \alpha_k = 0, \quad \beta^2 = 1 \quad (k, l = 1, 2, 3).$$

The operator  $D_{\text{hyd}}$  is a relativistic version of  $H_{\text{hyd}}$  [20].

It is shown [20, §7.4] that the discrete spectrum  $\sigma_{\text{disc}}(D_{\text{hyd}})$  of  $D_{\text{hyd}}$  is given by

$$\sigma_{\text{disc}}(D_{\text{hyd}}) = \{E_{n,j}\}_{n,j}$$

with

$$E_{n,j} = \frac{m_e c^2}{\sqrt{1 + \frac{1}{\hbar^2 c^2} \left( \frac{\gamma}{n - \left(j + \frac{1}{2}\right) + \sqrt{\left(j + \frac{1}{2}\right)^2 - \frac{\gamma^2}{\hbar^2 c^2}} \right)^2}}, \quad n = 1, 2, \dots,$$

where  $j$  ( $1/2 \leq j \leq n - 1/2$ ) is the total angular momentum of the electron, being related to the orbital angular momentum  $\ell = 0, 1, \dots$  by  $j = \ell \pm 1/2$  ( $\pm 1/2$  are the possible values of the spin of the electron), and the condition  $\gamma/\hbar c < 1$  is assumed.

It is easy to see that  $E_{n,j}$  is monotone increasing in  $n$  and that, for each  $n$ ,

$$E_{n,j} < E_{n,j+1}.$$

Note also that the non-relativistic limit<sup>2</sup>  $c \rightarrow \infty$  of  $E_{n,j} - m_e c^2$  gives  $E_n$ :

$$\lim_{c \rightarrow \infty} (E_{n,j} - m_e c^2) = E_n, \quad n = 1, 2, \dots.$$

For each  $n = 1, 2, \dots$ , the state with energy eigenvalue  $E_{n,j}$  and angular momentum  $\ell = 0, 1, 2, 3, 4, \dots$  is respectively labeled as  $nx_j$  with  $x = s$  ( $\ell = 0$ ),  $p$  ( $\ell = 1$ ),  $d$  ( $\ell = 2$ ),  $f$  ( $\ell = 3$ ),  $g$  ( $\ell = 4$ ),  $\dots$ :

principal number	state
$n = 1$	$1s_{1/2}$
$n = 2$	$(2s_{1/2}, 2p_{1/2}), \quad 2p_{3/2}$
$n = 3$	$(3s_{1/2}, 3p_{1/2}), \quad (3p_{3/2}, 3d_{3/2}), \quad 3d_{5/2}$
$\vdots$	$\vdots$

Here the states in each round bracket are degenerate. For example, the states  $2s_{1/2}$  and  $2p_{1/2}$  are degenerate with energy  $E_{2,1/2}$ . The energy levels  $E_{2,1/2}$  and  $E_{2,3/2}$  are very near with  $E_{2,1/2} < E_{2,3/2}$ . Hence these energy levels subtracted by  $m_e c^2$  may be regarded as a split of the second principal energy level  $E_2$  in the non-relativistic theory. It is known that the energy levels  $\{E_{n,j} - m_e c^2\}_{n,j}$  gives a good agreement with experimental data (Fig.1(b)).

In 1947, however, Lamb and Retherford [15] experimentally observed that there is a very small difference between the energies of the states  $2s_{1/2}$  and  $2p_{1/2}$  with the former being higher than the latter (Fig.2). This difference is called the *Lamb shift*. Thus the

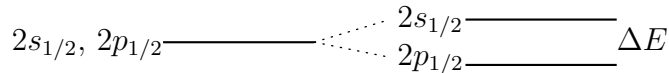


Figure 2:  $\Delta E$  =Lamb shift

Dirac theory breaks down in this respect.

It was Bethe [10] who first explained the Lamb shift using non-relativistic QED. He considered the Lamb shift as an energy shift caused by the interaction of the electron with the quantum radiation field. In his calculation, which is based on the standard heuristic perturbation theory, the mass renormalization of the electron is one of the essential

<sup>2</sup>In a non-relativistic theory, the kinetic energy of a rest particle is zero. Hence, in taking the non-relativistic limit of an energy in a relativistic theory, one must subtract the rest energy  $m_e c^2$  from it.

prescriptions. On the other hand, Welton [21] gave another method to explain the Lamb shift using non-relativistic QED: He infers that the interaction of the electron with the quantum radiation field may give rise to fluctuations of the position of the electron and these fluctuations may change the Coulomb potential so that the energy level shift such as the Lamb shift may occur. With this physical intuition, he derived the Lamb shift heuristically and perturbatively. After the work of Bethe and Welton, perturbative calculations of the Lamb shift using relativistic QED with prescription of renormalizations have been made, giving amazingly good agreements with the experimental result (see, e.g., [14]). However a mathematically rigorous construction of relativistic QED (existence of full relativistic QED) is still open as one of most important and challenging problems in modern mathematical physics. On the other hand, non-relativistic QED allows one to analyze it in a mathematically rigorous way [1, 2, 3](for a review of recent developments of non-relativistic QED, see, e.g., [13]).

Motivated by finding a mathematically general theory behind Welton's heuristic arguments made in [21], the present author developed in the paper [4] an abstract theory of scaling limit for self-adjoint operators on a Hilbert space and applied it to one-particle non-relativistic QED (a quantum mechanical model of a non-relativistic charged particle interacting with the quantum radiation field; a variant of the Pauli-Fierz model [16]) to obtain an effective Hamiltonian of the whole quantum system. This result is the starting point of the present review. Thus we next explain it in some detail.

## 2 A Model in Non-relativistic QED and Scaling Limit

For mathematical generality, the non-relativistic charged particle is assumed to appear in the  $d$ -dimensional Euclidean vector space  $\mathbb{R}^d$  with  $d \geq 2$ , so that the Hilbert space of state vectors for the charged particle is taken to be  $L^2(\mathbb{R}^d)$ . We consider the situation where the charged particle is under the influence of a scalar potential  $V : \mathbb{R}^d \rightarrow \mathbb{R}$  (Borel measurable). Then the non-relativistic Hamiltonian of the charged particle with mass  $m > 0$  is given by the Schrödinger operator

$$H(m) := -\frac{\hbar^2}{2m}\Delta + V. \quad (2.1)$$

On the other hand, the Hilbert space of state vectors of a photon is given by

$$\mathcal{H}_{\text{ph}} := \oplus^{d-1} L^2(\mathbb{R}^d),$$

the  $(d-1)$ -direct sum of  $L^2(\mathbb{R}^d)$ , where the number  $(d-1)$  in the present context means the freedom of polarization of a photon and  $\mathbb{R}^d$  here denotes the space of wave number vectors of a photon. Then the Hilbert space of state vectors for the quantum radiation field is given by the boson Fock space

$$\begin{aligned} \mathcal{F}_{\text{rad}} &:= \oplus_{n=0}^{\infty} \otimes_s^n \mathcal{H}_{\text{ph}} \\ &= \left\{ \Psi = \{\Psi^{(n)}\}_{n=0}^{\infty} \mid \Psi^{(n)} \in \otimes_s^n \mathcal{H}_{\text{ph}}, n \geq 0, \sum_{n=0}^{\infty} \|\Psi^{(n)}\|^2 < \infty \right\} \end{aligned}$$

over  $\mathcal{H}_{\text{ph}}$ , where  $\otimes_s^n \mathcal{H}_{\text{ph}}$  denotes the  $n$ -fold symmetric tensor product of  $\mathcal{H}_{\text{ph}}$  with  $\otimes_s^0 \mathcal{H}_{\text{ph}} := \mathbb{C}$  (the set of complex numbers) and  $\|\Psi^{(n)}\|$  denotes the norm of  $\Psi^{(n)}$ .

As is easily shown,  $\otimes_s^n \mathcal{H}_{\text{ph}}$  is identified with the Hilbert space of square integrable functions  $\psi^{(n)}((\mathbf{k}_1, s_1), (\mathbf{k}_2, s_2), \dots, (\mathbf{k}_n, s_n))$  on  $(\mathbb{R}^d \times \{1, \dots, d-1\})^n$  ( $\mathbf{k}_j \in \mathbb{R}^d, s_k \in \{1, \dots, d-1\}$ ) which are totally symmetric in the variables  $(\mathbf{k}_1, s_1), (\mathbf{k}_2, s_2), \dots, (\mathbf{k}_n, s_n)$ , where the isomorphism comes from the correspondence

$$S_n(\otimes_{j=1}^n \psi_j) \rightarrow \frac{1}{n!} \sum_{\sigma \in S_n} \psi_{\sigma(1)}(\mathbf{k}_1, s_1) \cdots \psi_{\sigma(n)}(\mathbf{k}_n, s_n), \quad \psi_j = (\psi_j(\cdot, s))_{s=1}^{d-1} \in \mathcal{H}_{\text{ph}}$$

with  $S_n$  being the symmetrization operator on  $\otimes^n \mathcal{H}_{\text{ph}}$  and  $S_n$  denotes the symmetry group of  $n$ -th order. We use this identification.

In the physical case  $d = 3$ , the energy of a photon with wave number vector  $\mathbf{k} \in \mathbb{R}^3$  is given by  $\hbar c|\mathbf{k}|$  (by Planck-Einstein-de Broglie relation,  $\hbar\mathbf{k}$  is the momentum of the photon with wave number vector  $\mathbf{k}$ ). Thus, in the case of general dimensions  $d$ , we assume that the energy of a photon with wave number vector  $\mathbf{k} \in \mathbb{R}^d$  is given by  $\hbar c\omega(\mathbf{k})$  with a function  $\omega : \mathbb{R}^d \rightarrow [0, \infty)$  such that  $0 < \omega(\mathbf{k}) < \infty$  for a.e. (almost everywhere)  $\mathbf{k} \in \mathbb{R}^d$  with respect to the Lebesgue measure on  $\mathbb{R}^d$ . Then the free Hamiltonian of the quantum radiation field is defined by

$$H_{\text{rad}} := \bigoplus_{n=0}^{\infty} \hbar c\omega^{(n)},$$

where  $\omega^{(0)} := 0$  and  $\omega^{(n)}$  is the multiplication operator by the function

$$\omega^{(n)}(\mathbf{k}_1, \dots, \mathbf{k}_n) := \sum_{j=1}^n \omega(\mathbf{k}_j)$$

on  $(\mathbb{R}^d \times \{1, \dots, d-1\})^n$ .

For each  $f \in \mathcal{H}_{\text{ph}}$ , there exists a densely defined closed linear operator  $a(f)$  on  $\mathcal{F}_{\text{rad}}$ , called the photon annihilation operator with test vector  $f$ , such that its adjoint  $a(f)^*$  takes the form

$$(a(f)^*\Psi)^{(0)} = 0, \quad (a(f)^*\Psi)^{(n)} = S_n(f \otimes \Psi^{(n-1)}), \quad \Psi = \{\Psi^{(n)}\}_{n=0}^{\infty} \in D(a(f)^*), n \geq 1$$

(for more details, see [5, Chapter 10]). The operators  $a(f)$  and  $a(g)^*$  ( $f, g \in \mathcal{H}_{\text{ph}}$ ) satisfy the commutation relations—canonical commutation relations (CCR)—

$$\begin{aligned} [a(f), a(g)^*] &= \langle f, g \rangle, \\ [a(f), a(g)] &= 0, \quad [a(f)^*, a(g)^*] = 0 \end{aligned}$$

on the subspace

$$\mathcal{F}_{\text{rad},0} := \{\Psi = \{\Psi^{(n)}\}_{n=0}^{\infty} \in \mathcal{F}_{\text{rad}} \mid \exists n_0 \text{ such that } \Psi^{(n)} = 0, \forall n \geq n_0\},$$

where  $[A, B] := AB - BA$  (commutator) and  $\langle \cdot, \cdot \rangle$  denotes inner product (anti-linear in the first variable, linear in the second one). Thus the set  $\{a(f), a(f)^* \mid f \in \mathcal{H}_{\text{ph}}\}$  gives a representation of the CCR indexed by  $\mathcal{H}_{\text{ph}}$ .

For a.e.  $\mathbf{k} \in \mathbb{R}^d$ , there exists an orthonormal system  $\{\mathbf{e}^{(s)}(\mathbf{k})\}_{s=1}^{d-1}$  of  $\mathbb{R}^d$  such that each vector  $\mathbf{e}^{(s)}(\mathbf{k}) = (e_1^{(s)}(\mathbf{k}), \dots, e_d^{(s)}(\mathbf{k}))$  is orthogonal to  $\mathbf{k}$ .

Let  $\rho$  be a real distribution on  $\mathbb{R}^d$  such that its Fourier transform  $\hat{\rho} = \hat{\rho}(\mathbf{k}) := (2\pi)^{-d/2} \int_{\mathbb{R}^d} e^{-i\mathbf{k}\cdot\mathbf{x}} \rho(\mathbf{x}) d\mathbf{x}$  (symbolic expression)—is a function satisfying

$$\frac{\hat{\rho}}{\omega^a} \in L^2(\mathbb{R}^d) \setminus \{0\}, \quad a = \frac{3}{2}, \frac{1}{2}.$$

Then the quantum radiation field  $A(\rho) := (A_1(\rho), \dots, A_d(\rho))$  smeared with  $\rho$  is defined by

$$A_j(\rho) = \frac{\sqrt{\hbar c}}{\sqrt{2}} \left\{ a \left( \frac{\hat{\rho}}{\sqrt{\omega}} e_j \right)^* + a \left( \frac{\hat{\rho}}{\sqrt{\omega}} e_j \right) \right\}, \quad j = 1, \dots, d,$$

where  $e_j : \mathbb{R}^d \rightarrow \mathbb{R}^{d-1}$ ,  $e_j(\mathbf{k}) := (e_j^{(1)}(\mathbf{k}), \dots, e_j^{(d-1)}(\mathbf{k}))$ , a.e.  $\mathbf{k} \in \mathbb{R}^d$ . We remark that, for the definition of  $A_j(\rho)$  itself, condition  $\hat{\rho}/\sqrt{\omega} \in L^2(\mathbb{R}^d)$  is sufficient. The additional condition  $\hat{\rho}/\omega^{3/2} \in L^2(\mathbb{R}^d)$  is needed in the development below.

The Hilbert space  $\mathfrak{H}$  of state vectors of the quantum system under consideration is given by

$$\mathfrak{H} = L^2(\mathbb{R}^d) \otimes \mathcal{F}_{\text{rad}}.$$

The Hamiltonian of our model is of the following form:

$$H_{\text{NR}} = H(m_0) \otimes I + I \otimes H_{\text{rad}} + H_{\text{I}}(\rho, m_0)$$

with

$$H_{\text{I}}(\rho, m_0) := -\frac{q}{m_0 c} \sum_{j=1}^d p_j \otimes A_j(\rho),$$

where  $m_0 > 0$  is the “bare” mass of the particle (the mass of the particle before going into the interaction with the quantum radiation field),  $q \in \mathbb{R}$  and  $p_j := -i\hbar D_j$  denote respectively the electric charge and the momentum operator of the particle. The operator  $H_{\text{I}}(\rho, m_0)$  describes an interaction of the charged particle with the quantum radiation field. In this context, the function  $\hat{\rho}$  plays a role of momentum cutoff for photons interacting with the particle.

To draw from the Hamiltonian  $H_{\text{NR}}$  observable effects that the quantum field may give rise to the quantum particle, we consider a scaling limit of  $H_{\text{NR}}$ . Thus we introduce the following scaled Hamiltonian:

$$H_{\text{NR}}(\kappa) := H(m(\kappa)) \otimes I + \kappa I \otimes H_{\text{rad}} + \kappa H_{\text{I}}(\rho, m), \quad \kappa > 0,$$

with

$$\frac{1}{m(\kappa)} := \frac{1}{m} + \kappa \frac{(d-1)}{d} \left( \frac{q}{mc} \right)^2 \int_{\mathbb{R}^d} \frac{|\hat{\rho}(\mathbf{k})|^2}{\omega(\mathbf{k})^2} d\mathbf{k},$$

where  $m > 0$  is the observed mass of the particle.

Under the assumption that  $V$  is infinitesimally small with respect to  $-\Delta$ , the operator  $H_{\text{NR}}(\kappa)$  is self-adjoint and bounded below [4, Lemma 3.1].

**Remark 2.1** The scaled Hamiltonian  $H_{\text{NR}}(\kappa)$  is obtained by the scaling  $c \rightarrow \kappa c$  and  $q \rightarrow \kappa^{3/2} q$  in  $H_{\text{NR}}$  with  $H(m_0)$  and  $H_{\text{I}}(\rho, m_0)$  replaced by  $H(m(1))$  and  $H_{\text{I}}(\rho, m)$  respectively.



Replacing  $m_0$  with  $m(\kappa)$  is called a mass renormalization<sup>3</sup>. We want to emphasize that the mass renormalization makes the Hamiltonian bounded below (under the condition that  $H(m)$  is bounded below) [4, Lemma 3.1].

A scaling limit of the original Pauli-Fierz model with dipole approximation is discussed in [12] (see also [13]).

The vector

$$\Omega_0 := \{1, 0, 0, \dots\} \in \mathcal{F}_{\text{rad}} \quad (\Omega^{(0)} = 1, \Omega^{(n)} = 0, n \geq 1)$$

is called the Fock vacuum in  $\mathcal{F}_{\text{rad}}$ . We denote by  $P_0$  the orthogonal projection onto the 1-dimensional subspace  $\{\alpha\Omega_0 | \alpha \in \mathbb{C}\}$  spanned by  $\Omega_0$ .

It is shown that the operator

$$T := \frac{iq}{mc} \sum_{j=1}^d p_j \otimes \frac{1}{\sqrt{2\hbar c}} \left\{ a \left( \frac{\hat{\rho}}{\omega^{3/2}} e_j \right)^* - a \left( \frac{\hat{\rho}}{\omega^{3/2}} e_j \right) \right\}$$

is essentially self-adjoint. We denote its closure by  $\bar{T}$ .

Let

$$\lambda_q := \frac{(d-1)}{4d} \left( \frac{\hbar}{mc} \right)^2 \frac{q^2}{\hbar c} \int_{\mathbb{R}^d} \frac{|\hat{\rho}(\mathbf{k})|^2}{\omega(\mathbf{k})^3} d\mathbf{k}.$$

The following theorem is proved [4, Theorem 3.4]:

**Theorem 2.2** *Assume that  $\hat{\rho}$  is spherically symmetric (i.e., it depends only on  $|\mathbf{k}|$ ). Suppose that  $V$  satisfies the following two conditions:*

(V.1)  $D(\Delta) \subset D(V)$  and, for all  $a > 0$ ,  $V(-\Delta + a)^{-1}$  is bounded with  $\lim_{a \rightarrow \infty} \|V(-\Delta + a)^{-1}\| = 0$ .

(V.2) For all  $t > 0$ ,  $\int_{\mathbb{R}^d} e^{-t|\mathbf{y}|^2} |V(\mathbf{y})| d\mathbf{y} < \infty$ .

Then, for all  $z \in \mathbb{C} \setminus \mathbb{R}$ ,

$$s\text{-}\lim_{\kappa \rightarrow \infty} (H_{\text{NR}}(\kappa) - z)^{-1} = e^{-i\bar{T}} \left( (H_{\text{eff}} - z)^{-1} \otimes P_0 \right) e^{i\bar{T}},$$

where  $s\text{-}\lim$  means strong limit and

$$H_{\text{eff}} := -\frac{\hbar^2}{2m} \Delta + V_{\text{eff}} \tag{2.2}$$

with

$$V_{\text{eff}}(\mathbf{x}) := \frac{1}{(4\pi\lambda_q)^{d/2}} \int_{\mathbb{R}^d} e^{-|\mathbf{x}-\mathbf{y}|^2/4\lambda_q} V(\mathbf{y}) d\mathbf{y}, \quad \mathbf{x} \in \mathbb{R}^d. \tag{2.3}$$

---

<sup>3</sup>Strictly speaking, one should replace  $m_0$  in  $H_{\text{I}}(\rho, m_0)$  with  $m(\kappa)$  too. But, since

$$H_{\text{I}}(\rho, m(\kappa)) = H_{\text{I}}(\rho, m) - \kappa \frac{q}{c} \frac{(d-1)}{d} \left( \frac{q}{mc} \right)^2 \left( \int_{\mathbb{R}^d} \frac{|\hat{\rho}(\mathbf{k})|^2}{\omega(\mathbf{k})^2} d\mathbf{k} \right) \sum_{j=1}^d p_j \otimes A_j(\rho).$$

and the second term on the right hand side is of the third order in  $q$ , one may take into account only the first term on the right hand side as a primary approximation in a perturbative sense.

**Remark 2.3** Under condition (V.1),  $V$  is infinitesimally small with respect to  $-\Delta$  and hence  $H(m_0)$  is self-adjoint and bounded below for all  $m_0 > 0$ . Moreover, under conditions (V.1) and (V.2),  $V_{\text{eff}}$  is infinitesimally small with respect to  $-\Delta$  and hence  $H_{\text{eff}}$  is self-adjoint and bounded below (see [4, §III, B]).

Theorem 2.2 may be physically interpreted as follows: the limiting system as  $\kappa \rightarrow \infty$  restricted to the subspace  $L^2(\mathbb{R}^3) \otimes \{\alpha e^{-i\bar{T}} \Omega_0 | \alpha \in \mathbb{C}\}$  is equivalent to the particle system whose Hamiltonian is  $H_{\text{eff}}$ . Therefore  $H_{\text{eff}}$  may include observable effects of the original interacting system through  $V_{\text{eff}}$ . In this sense, we call  $V_{\text{eff}}$  an effective potential for the particle system and, correspondingly to this, we call  $H_{\text{eff}}$  an effective Hamiltonian of the particle interacting with the quantum radiation field.

To see if the effective Hamiltonian  $H_{\text{eff}}$  really explains some observable effects, one has to investigate the spectral properties of it. This was the main motivation of the paper [7]. In what follows, we concentrate our attention on this aspect.

### 3 Elementary Properties of the Effective Hamiltonian

It is obvious that  $q \rightarrow 0$  if and only if  $\lambda_q \rightarrow 0$ . Hence we replace  $\lambda_q$  by a parameter  $\lambda > 0$  and regard  $\lambda$  as a perturbation parameter, where the limit  $\lambda \downarrow 0$  corresponds to the unperturbed case. Thus we consider the effective Hamiltonian in the form

$$H_\lambda := -\frac{\hbar^2}{2m}\Delta + V_\lambda, \quad \lambda > 0, \quad (3.1)$$

with

$$V_\lambda(\mathbf{x}) := \frac{1}{(4\pi\lambda)^{d/2}} \int_{\mathbb{R}^d} e^{-|\mathbf{x}-\mathbf{y}|^2/4\lambda} V(\mathbf{y}) d\mathbf{y}. \quad (3.2)$$

As for  $V$ , we assume only that

$$\int_{\mathbb{R}^d} e^{-t|\mathbf{y}|^2} |V(\mathbf{y})| d\mathbf{y} < \infty, \quad \forall t > 0. \quad (3.3)$$

We have

$$H_{\text{eff}} = H_{\lambda_q}. \quad (3.4)$$

**Remark 3.1** If  $V \in L^p(\mathbb{R}^d)$  for some  $1 \leq p \leq \infty$ , then (3.3) is satisfied.

Note that  $V_\lambda$  is the convolution of  $V$  and the Gaussian function

$$G_\lambda(\mathbf{x}) := \frac{1}{(4\pi\lambda)^{d/2}} e^{-\mathbf{x}^2/4\lambda}, \quad \mathbf{x} \in \mathbb{R}^d, \quad (3.5)$$

i.e.,

$$V_\lambda = G_\lambda * V. \quad (3.6)$$

In other words,  $V_\lambda$  is the Gauss transform of  $V$  with the Gaussian function  $G_\lambda$ . This structure may be suggestive, because the function  $G_\lambda(\mathbf{x} - \mathbf{y})$  of  $\mathbf{x}$  and  $\mathbf{y}$  is the integral kernel of the heat semi-group  $\{e^{\lambda\Delta}\}_{\lambda>0}$  on  $L^2(\mathbb{R}^d)$  (the heat kernel).

The effective potential  $V_\lambda$  is a perturbation of  $V$  in the following senses:

- (i) If  $V$  is continuous and  $\sup_{\mathbf{x} \in \mathbb{R}^d} |V(\mathbf{x})| e^{-c|\mathbf{x}|^\alpha} < \infty$  for some  $c > 0$  and  $\alpha \in [0, 2)$ , then

$$\lim_{\lambda \downarrow 0} V_\lambda(\mathbf{x}) = V(\mathbf{x}), \quad \mathbf{x} \in \mathbb{R}^d.$$

- (ii) If  $V \in L^2(\mathbb{R}^d)$ , then

$$V_\lambda = e^{\lambda \Delta} V \in L^2(\mathbb{R}^d).$$

and hence  $\lim_{\lambda \rightarrow 0} \|V_\lambda - V\|_{L^2(\mathbb{R}^d)} = 0$  holds<sup>4</sup>.

- (iii) If  $V \in L^p(\mathbb{R}^d)$  for some  $p \in [1, \infty)$ , then  $V_\lambda \in L^p(\mathbb{R}^d)$  and

$$\lim_{\lambda \rightarrow 0} \|V_\lambda - V\|_{L^p(\mathbb{R}^d)} = 0.$$

- (iv) If  $V \in L^\infty(\mathbb{R}^d)$  and  $V$  is uniformly continuous on  $\mathbb{R}^d$ , then  $V_\lambda \in L^\infty(\mathbb{R}^d)$  and

$$\lim_{\lambda \rightarrow 0} \|V_\lambda - V\|_{L^\infty(\mathbb{R}^d)} = 0.$$

Thus, from a perturbation theoretical point of view, it is natural to write

$$H_\lambda = H_0 + W_\lambda \tag{3.7}$$

with

$$H_0 := H(m) = -\frac{\hbar^2}{2m} \Delta + V, \tag{3.8}$$

$$W_\lambda := V_\lambda - V. \tag{3.9}$$

However, we want to emphasize that  $H_\lambda$  is not necessarily a regular perturbation of  $H$  in the sense of [19, §XII.2]. Even in that case, the order of the perturbation may be infinite.

One can analyze general aspects of spectra of  $H_\lambda$  [7]. But, here, we restrict ourselves to the case where  $V$  is a spherically symmetric function on  $\mathbb{R}^3$ .

## 4 Spectral Properties of $H_\lambda$ with a Spherically Symmetric Potential $V$ on $\mathbb{R}^3$

We consider the case where  $d = 3$  and  $V$  is given by the following form:

$$V(\mathbf{x}) = \frac{u(|\mathbf{x}|)}{|\mathbf{x}|}, \quad \mathbf{x} \in \mathbb{R}^3 \setminus \{\mathbf{0}\} \tag{4.1}$$

with  $u : [0, \infty) \rightarrow \mathbb{R}$  being *bounded and continuously differentiable on  $[0, \infty)$  with the derivative  $u'$  bounded on  $[0, \infty)$* . Note that  $V$  has singularity at  $\mathbf{x} = \mathbf{0}$  if  $u(\mathbf{0}) \neq 0$ . It is

---

<sup>4</sup>For  $p \in [1, \infty]$ ,  $\|\cdot\|_{L^p(\mathbb{R}^d)}$  denotes the norm of  $L^p(\mathbb{R}^d)$ .

easy to see that this  $V$  satisfies condition (3.3). By direct computations, one sees that the effective potential  $V_\lambda$  in the present case takes the form

$$V_\lambda(\mathbf{x}) = \frac{e^{-|\mathbf{x}|^2/4\lambda}}{\sqrt{\pi\lambda}|\mathbf{x}|} \int_0^\infty e^{-r^2/4\lambda} u(r) \sinh \frac{|\mathbf{x}|r}{2\lambda} dr. \quad (4.2)$$

In particular,  $V_\lambda$  also is spherically symmetric<sup>5</sup>.

A basic result on the spectra of  $H_\lambda$  is stated in the next theorem:

**Theorem 4.1** *Let  $V$  be given by (4.1). Then, for all  $\lambda \geq 0$ ,  $H_\lambda$  is self-adjoint with  $D(H_\lambda) = D(\Delta)$  and bounded below. Moreover*

$$\sigma_{\text{ess}}(H_\lambda) = [0, \infty),$$

where  $\sigma_{\text{ess}}(\cdot)$  denotes essential spectrum, and, if there exists an  $r_0 > 0$  such that  $\sup_{r \geq r_0} u(r) < 0$ , then the discrete spectrum  $\sigma_{\text{disc}}(H_\lambda)$  is infinite.

Suppose that  $H_0$  has an isolated eigenvalue  $E_0 \in \mathbb{R}$  with finite multiplicity  $m(E_0)$  ( $1 \leq m(E_0) < \infty$ ). Let  $r$  be a constant satisfying

$$0 < r < \min_{E \in \sigma(H_0) \setminus \{E_0\}} |E - E_0|.$$

Then

$$C_r(E_0) := \{z \in \mathbb{C} \mid |z - E_0| = r\} \subset \rho(H_0),$$

Let

$$n_r := r \sup_{z \in C_r(E_0)} \|(H_0 - z)^{-1}\|, \quad r_\lambda := \sup_{z \in C_r(E_0)} \|W_\lambda(H_0 - z)^{-1}\|.$$

**Theorem 4.2** *Let  $r_\lambda < 1/(1 + n_r)$ . Then,  $H_\lambda$  has exactly  $m(E_0)$  eigenvalues in the interval  $(E_0 - r, E_0 + r)$ , counting multiplicities, and  $\sigma(H_\lambda) \cap (E_0 - r, E_0 + r)$  consists of only these eigenvalues.*

In the case where  $E_0$  is a simple eigenvalue of  $H$ , one can obtain more detailed results:

**Corollary 4.3** *Let  $r_\lambda < 1/(1 + n_r)$ . Suppose that  $m(E_0) = 1$  and  $\Omega_0$  is a normalized eigenvector of  $H$  with eigenvalue  $E_0$ . Then,  $H_\lambda$  has exactly one simple eigenvalue  $E_\lambda$  in the interval  $(E_0 - r, E_0 + r)$  with formula*

$$E_\lambda = E_0 + \frac{\langle \Omega_0, W_\lambda \Omega_0 \rangle + \sum_{n=1}^\infty S_n(\lambda)}{1 + \sum_{n=1}^\infty T_n(\lambda)},$$

where

$$S_n(\lambda) := \frac{(-1)^{n+1}}{2\pi i} \int_{C_r(E_0)} dz \left\langle \Omega_0, [W_\lambda(H - z)^{-1}]^{n+1} \Omega_0 \right\rangle,$$

$$T_n(\lambda) := \frac{(-1)^{n+1}}{2\pi i} \int_{C_r(E_0)} dz \frac{\langle \Omega_0, [W_\lambda(H - z)^{-1}]^n \Omega_0 \rangle}{E_0 - z},$$

<sup>5</sup>It is an easy exercise to show that, if  $V$  is spherically symmetric on  $\mathbb{R}^d$ , then so is  $V_\lambda$ .

and  $\sigma(H_\lambda) \cap (E_0 - r, E_0 + r) = \{E_\lambda\}$ . Moreover, a normalized eigenvector of  $H_\lambda$  with eigenvalue  $E_\lambda$  is given by

$$\Omega_\lambda = \frac{\Omega_0 + \sum_{n=1}^{\infty} \Omega_{\lambda,n}}{\sqrt{1 + \sum_{n=1}^{\infty} T_n(\lambda)}},$$

where

$$\Omega_{\lambda,n} := \frac{(-1)^{n+1}}{2\pi i} \int_{C_r(E_0)} dz (H - z)^{-1} [W_\lambda(H - z)^{-1}]^n \Omega_0.$$

## 5 Reductions of $H_\lambda$ to Closed Subspaces

The Hilbert space  $L^2(\mathbb{R}^3)$  has the orthogonal decomposition

$$L^2(\mathbb{R}^3) = \bigoplus_{\ell=0}^{\infty} \bigoplus_{s=-\ell}^{\ell} \mathcal{H}_\ell^s$$

with

$$\mathcal{H}_\ell^s = L^2([0, \infty), r^2 dr) \otimes \{\alpha Y_\ell^s | \alpha \in \mathbb{C}\},$$

where  $Y_\ell^s$  is the spherical harmonics with index  $(\ell, s)$ :

$$Y_\ell^s(\theta, \phi) := (-1)^s \sqrt{\frac{(\ell-s)!}{(\ell+s)!}} \sqrt{\frac{2\ell+1}{4\pi}} P_\ell^s(\cos \theta) e^{is\phi},$$

$$\theta \in [0, \pi], \phi \in [0, 2\pi), s = -\ell, -\ell+1, \dots, 0, \dots, \ell-1, \ell$$

with  $P_\ell^s$  being the associated Legendre function:

$$P_\ell^s(x) := (1-x^2)^{s/2} \frac{d^s}{dx^s} \frac{(-1)^\ell}{2^\ell \ell!} \left( \frac{d}{dx} \right)^\ell (1-x^2)^\ell, \quad |x| < 1.$$

We have

$$\int_0^\pi d\theta \int_0^{2\pi} d\phi \sin \theta Y_\ell^s(\theta, \phi) {}^* Y_{\ell'}^{s'}(\theta, \phi) = \delta_{\ell\ell'} \delta_{ss'}.$$

As we have already seen,  $V_\lambda$  under consideration is spherically symmetric. Hence  $H_\lambda$  is reduced by each  $\mathcal{H}_\ell^s$ . We denote the reduced part of  $H_\lambda$  by  $H_\lambda^{\ell,s}$ :

$$(H_\lambda^{\ell,s} f \otimes Y_\ell^s)(r, \phi, \theta) = \left( -\frac{\hbar^2}{2m} \frac{d^2}{dr^2} + \tilde{V}_\lambda(r) - \frac{\hbar^2}{2m} \frac{2}{r} \frac{d}{dr} \right) f(r) Y_\ell^s(\theta, \phi)$$

$$+ \frac{\ell(\ell+1)}{r^2} f(r) Y_\ell^s(\theta, \phi), \quad f \in C_0^\infty(0, \infty),$$

where  $\tilde{V}_\lambda(r) := V_\lambda(\mathbf{x})|_{r=|\mathbf{x}|}$  and  $C_0^\infty(0, \infty)$  is the set of infinitely differentiable functions on  $(0, \infty)$  with bounded support in  $(0, \infty)$ .

**Corollary 5.1** *For each pair  $(\ell, s)$  ( $\ell \in \{0\} \cup \mathbb{N}, s = -\ell, -\ell+1, \dots, \ell$ ), Theorem 4.2 and Corollary 4.3 with  $H_\lambda$  replaced by  $H_\lambda^{\ell,s}$  hold.*

## 6 Energy Level Shifts in a Hydrogen-like Atom

Now we consider a hydrogen-like atom mentioned in Introduction. Thus we take as an unperturbed Hamiltonian  $H_0$  the Schrödinger operator  $H_{\text{hyd}}$  defined by (1.1):

$$H_{\text{hyd}} = -\frac{\hbar^2}{2m_e}\Delta + V^{(\gamma)}, \quad V^{(\gamma)} := -\frac{\gamma}{|\mathbf{x}|}. \quad (6.1)$$

The eigenvalue  $E_n$  of  $H_{\text{hyd}}$  (see (1.3)) is a unique simple eigenvalue of the reduced part  $H_{\text{hyd}}^{\ell,s}$  of  $H_{\text{hyd}}$  ( $0 \leq \ell \leq n-1$ ) to the closed subspace  $\mathcal{H}_\ell^s$  with a normalized eigenfunction

$$\begin{aligned} \psi_{n,\ell,s}(\mathbf{x}) &:= C_{n,\ell} e^{-\beta_n r/2} (\beta_n r)^\ell L_{n+\ell}^{2\ell+1}(\beta_n r) Y_\ell^s(\theta, \phi), \\ r &= |\mathbf{x}|, \ell = 0, 1, \dots, n-1, \end{aligned}$$

where

$$\beta_n := \frac{2m_e \gamma}{\hbar^2 n},$$

$L_n^k$  ( $0 \leq k \leq n$ ) is the Laguerre associated polynomial with order  $n-k$ , i.e.,

$$L_n^k(x) = \frac{d^k}{dx^k} L_n(x), \quad x \in \mathbb{R}$$

with  $L_n(x)$  being the  $n$ -th Laguerre polynomial and

$$C_{n,\ell} := \frac{\beta_n^{3/2} \sqrt{(n-\ell-1)!}}{\sqrt{[(n+\ell)!]^3 2n}}.$$

Applying (4.2) with  $u = -\gamma$  (a constant function), the effective potential

$$V_\lambda^{(\gamma)} := G_\lambda * V^{(\gamma)}$$

in the present case is of the form:

$$V_\lambda^{(\gamma)} = V^{(\gamma)} + W_\lambda^{(\gamma)}$$

with

$$W_\lambda^{(\gamma)}(\mathbf{x}) := \frac{2\gamma}{\sqrt{\pi}|\mathbf{x}|} \text{Erfc}(|\mathbf{x}|/2\sqrt{\lambda}),$$

where  $\text{Erfc} : \mathbb{R} \rightarrow [0, \infty)$  is the Gauss error function:

$$\text{Erfc}(x) := \int_x^\infty e^{-y^2} dy, \quad x \geq 0.$$

Hence the effective Hamiltonian

$$H_\lambda(\gamma) = -\frac{\hbar^2}{2m_e}\Delta + V_\lambda^{(\gamma)}, \quad \lambda > 0,$$

takes the form

$$H_\lambda(\gamma) = H_{\text{hyd}} + W_\lambda^{(\gamma)}.$$

The next theorem follows from a simple application of Theorem 4.1:

**Theorem 6.1** For all  $\lambda > 0$  and  $\gamma > 0$ ,  $H_\lambda(\gamma)$  is self-adjoint with  $D(H_\lambda(\gamma)) = D(\Delta)$  and bounded below. Moreover,  $\sigma_{\text{disc}}(H_\lambda(\gamma))$  is infinite and

$$\sigma_{\text{disc}}(H_\lambda(\gamma)) \subset (-\infty, 0), \quad \sigma_{\text{ess}}(H_\lambda(\gamma)) = [0, \infty).$$

We take  $r_n > 0$  such that  $r_n < |E_{n+1} - E_n|$  and set

$$C_{r_n}(E_n) := \{z \in \mathbb{C} \mid |z - E_n| = r_n\}.$$

Let

$$M_n := r_n \sup_{z \in C_{r_n}(E_n)} \|H(\gamma) - z\|^{-1}, \quad R_{\lambda,n} := \sup_{z \in C_{r_n}(E_n)} \|W_\lambda^{(\gamma)}(H(\gamma) - z)^{-1}\|.$$

We denote by  $H_\lambda^{\ell,s}(\gamma)$  the reduced part of  $H_\lambda(\gamma)$  to  $\mathcal{H}_\ell^s$ .

We have from Corollary 5.1 the following result:

**Theorem 6.2** Let  $n \in \mathbb{N}$ ,  $\ell = 0, 1, \dots, n-1$  and  $s = -\ell, -\ell+1, \dots, \ell$ . Suppose that  $\lambda > 0$  and  $R_{\lambda,n} < 1/(1 + M_n)$ . Then,  $H_\lambda^{\ell,s}(\gamma)$  has a unique simple eigenvalue  $E_{n,\ell,s}(\lambda)$  near  $E_n$  with

$$E_{n,\ell,s}(\lambda) = E_n + \frac{\left\langle \psi_{n,\ell,s}, W_\lambda^{(\gamma)} \psi_{n,\ell,s} \right\rangle + \sum_{p=1}^{\infty} F_{n,\ell,s}^{(p)}(\lambda)}{1 + \sum_{p=1}^{\infty} G_{n,\ell,s}^{(p)}(\lambda)},$$

where

$$F_{n,\ell,s}^{(p)}(\lambda) := \frac{(-1)^{p+1}}{2\pi i} \int_{C_{r_n}(E_n)} \left\langle \psi_{n,\ell,s}, \left[ W_\lambda^{(\gamma)}(H(\gamma) - z)^{-1} \right]^{p+1} \psi_{n,\ell,s} \right\rangle dz,$$

$$G_{n,\ell,s}^{(p)}(\lambda) := \frac{(-1)^{p+1}}{2\pi i} \int_{C_{r_n}(E_n)} \frac{\left\langle \psi_{n,\ell,s}, \left[ W_\lambda^{(\gamma)}(H(\gamma) - z)^{-1} \right]^p \psi_{n,\ell,s} \right\rangle}{E_n - z} dz.$$

Moreover, a normalized eigenvector  $\psi_{n,\ell,s}^{(\lambda)}$  of  $H_\lambda^{\ell,s}(\gamma)$  with eigenvalue  $E_{n,\ell,s}(\lambda)$  is given by

$$\psi_{n,\ell,s}^{(\lambda)} = \frac{\psi_{n,\ell,s} + \sum_{p=1}^{\infty} S_{n,\ell,s}^{(p)}(\lambda)}{\sqrt{1 + \sum_{p=1}^{\infty} G_{n,\ell,s}^{(p)}(\lambda)}},$$

where

$$S_{n,\ell,s}^{(p)}(\lambda) := \frac{(-1)^{p+1}}{2\pi i} \int_{C_{r_n}(E_n)} (H(\gamma) - z)^{-1} \left[ W_\lambda^{(\gamma)}(H(\gamma) - z)^{-1} \right]^p \psi_{n,\ell,s} dz.$$

Let  $n \in \mathbb{N}$ ,  $\lambda > 0$  and  $R_{\lambda,n} < 1/(1 + M_n)$ . Then, by Theorem 6.2, one can define

$$\Delta E_n(\ell, s; \ell', s') := E_{n,\ell,s}(\lambda) - E_{n,\ell',s'}(\lambda) \quad (6.2)$$

for  $\ell, \ell' = 0, 1, \dots, n-1$ ,  $s, s' = -\ell, -\ell+1, \dots, \ell$  with  $(\ell, s) \neq (\ell', s')$ . We call it an energy level shift of  $H_\lambda(\gamma)$  with respect to the  $n$ -th energy level.

The next theorem is an important result necessary for deriving the Lamb shift (see the next section):

**Theorem 6.3** Under the assumption of Theorem 6.2, the following holds:

$$E_{n,\ell,s}(\lambda) = E_n + 4\pi\gamma |\psi_{n,\ell,s}(\mathbf{0})|^2 \lambda + o(\lambda) \quad (\lambda \rightarrow 0). \quad (6.3)$$

## 7 Derivation of the Lamb shift

In this section, we assume that, for each  $n \in \mathbb{N}$ ,  $\lambda > 0$  is sufficiently small so that the assumption of Theorem 6.2 holds. Then, by Theorem 6.3, we have

$$\Delta E_n(\ell, s; \ell', s') = 4\pi\gamma(|\psi_{n,\ell,s}(\mathbf{0})|^2 - |\psi_{n,\ell',s'}(\mathbf{0})|^2)\lambda + o(\lambda) \quad (\lambda \rightarrow 0).$$

Using

$$L_n^1(0) = nn!, \quad Y_0^0 = \frac{1}{\sqrt{4\pi}},$$

we obtain

$$|\psi_{n,\ell,s}(\mathbf{0})|^2 = \begin{cases} \frac{1}{\pi} \left( \frac{m_e \gamma}{\hbar^2} \right)^3 \frac{1}{n^3} & ; \ell = 0, s = 0 \\ 0 & ; \ell \geq 1 \end{cases} \quad (7.1)$$

Hence the following hold:

(i) If  $\ell, \ell' \geq 1$ , then

$$\Delta E_n(\ell, s; \ell', s') = o(\lambda) \quad (\lambda \rightarrow 0). \quad (7.2)$$

(ii) If  $\ell \geq 1$ , then

$$\Delta E_n(0, 0; \ell, s) = 4\pi\gamma\lambda|\psi_{n,0,0}(\mathbf{0})|^2 + o(\lambda) \quad (\lambda \rightarrow 0). \quad (7.3)$$

Formula (7.3) shows that, for each  $n$ , the energy of the state with  $\ell = 0, s = 0$  (the  $s$ -state) is higher than that of the state with  $\ell \geq 1$  for all sufficiently small  $\lambda$ . This may be a non-relativistic correspondence of the experimental fact that, for  $n = 2$ , the energy of the state  $2s_{1/2}$  is higher than that of the state  $2p_{1/2}$ .

To compare the value of  $\Delta E_n(0, 0; \ell, s)$  with the experimental one, we take  $\lambda = \lambda_q$  with  $q = -e$ ,  $m = m_e$  and

$$\omega(\mathbf{k}) = |\mathbf{k}|, \quad \hat{\rho}(\mathbf{k}) = \frac{1}{\sqrt{(2\pi)^3}} \chi_{[\omega_{\min}/\hbar c, \omega_{\max}/\hbar c]}(|\mathbf{k}|), \quad \mathbf{k} \in \mathbb{R}^3,$$

with constants  $\omega_{\min} > 0$  and  $\omega_{\max} > 0$  satisfying  $\omega_{\min} < \omega_{\max}$ . Then we have

$$\lambda = \lambda_{-e} = \alpha \left( \frac{\hbar}{m_e c} \right)^2 \frac{1}{3\pi} \log \frac{\omega_{\max}}{\omega_{\min}},$$

where

$$\alpha := \frac{e^2}{4\pi\hbar c} \approx \frac{1}{137}$$

is the fine structure constant. We remark that  $\omega_{\min}$  (resp.  $\omega_{\max}$ ) physically means an infrared (resp. ultraviolet) cutoff of the one-photon energy. We have  $\gamma = Ze^2/4\pi$ . Thus we obtain

$$\begin{aligned} \Delta E_n(0, 0; \ell, s) &\approx \alpha^5 \frac{4}{3\pi} m_e c^2 \frac{Z^4}{n^3} \log \frac{\omega_{\max}}{\omega_{\min}} \\ &= \frac{8}{3\pi} \alpha^3 \text{Ry} \frac{Z^4}{n^3} \log \frac{\omega_{\max}}{\omega_{\min}} \quad (\alpha \rightarrow 0), \end{aligned} \quad (7.4)$$



where  $\text{Ry} := \alpha^2 m_e c^2 / 2$  is 1 rydberg ( $-\text{Ry}$  is the ground state energy of the hydrogen atom). If we take  $\omega_{\max} = m_e c^2$  (the rest mass energy of the electron) and  $\omega_{\min} = 17.8 \text{ Ry}$ , then the right hand side of (7.4) completely coincides with Bethe's calculation [10] of the Lamb shift. Hence it is in a good agreement with the experimental result.

**Remark 7.1** There is another method to estimate the energy level shifts of atoms due to the interaction with the quantum radiation field [11], using the Pauli-Fierz model with spin in non-relativistic QED.

## 8 A General Class of Abstract Effective Operators

In concluding this paper, we want to point out that the effective Hamiltonian  $H_\lambda$  given by (3.1) is a special case of an abstract effective operator obtained as a scaling limit of the generalized spin-boson (GSB) model [8], a general model of an abstract “particle” quantum system coupled to a Bose field (a quantum field of bosons).

The GSB model is described as follows. The Hilbert space for the abstract “particle” quantum system is taken to be an abstract complex Hilbert space  $\mathcal{H}$  and the Hilbert space for the Bose field is given by the boson Fock space

$$\mathcal{F}_b(\mathcal{K}) := \bigoplus_{n=0}^{\infty} \otimes_s^n \mathcal{K}$$

over an abstract Hilbert space  $\mathcal{K}$  ( $\mathcal{F}_b(\mathcal{K})$  is just the Hilbert space  $\mathcal{F}_{\text{rad}}$  with  $\mathcal{H}_{\text{ph}}$  replaced with  $\mathcal{K}$ ). Then the Hilbert space of the composite system is the tensor product  $\mathcal{H} \otimes \mathcal{F}_b(\mathcal{K})$ .

The abstract version of the photon annihilation operator on  $\mathcal{F}_{\text{rad}}$  is defined on  $\mathcal{F}_b(\mathcal{K})$  too, in exactly the same form. We denote it by  $a(f)$ ,  $f \in \mathcal{K}$  and set

$$\phi(f) := \frac{1}{\sqrt{2}}(a(f)^* + a(f)), \quad f \in \mathcal{K},$$

the Segal field operator, which is essentially self-adjoint ([5, §4.6], [18, §X.7]).

Let  $T$  be a non-negative, injective self-adjoint operator on  $\mathcal{K}$  denoting the one-particle Hamiltonian of the Bose field. Then the  $n$ -particle Hamiltonian ( $n \geq 0$ ) of the Bose field is defined by  $T^{(0)} := 0$  and

$$T^{(n)} := \sum_{j=1}^n I \otimes \cdots \otimes I \otimes \overbrace{T}^{j\text{th}} \otimes I \cdots \otimes I, \quad n \geq 1.$$

The free Hamiltonian of the Bose field is given by

$$H_b = d\Gamma(T) := \bigoplus_{n=0}^{\infty} T^{(n)},$$

the second quantization of  $T$  ([5, §4.4.2], [17, p.302]).

Let  $A$  be a self-adjoint operator on  $\mathcal{H}$  denoting the Hamiltonian of the particle system,  $B_j$  ( $j = 1, \dots, N$ ) be a symmetric operator on  $\mathcal{H}$ , and  $g_j \in \mathcal{K}$  ( $j = 1, \dots, N$ ). Then the Hamiltonian of the GSB model is defined by

$$H_{\text{GSB}} := A \otimes I + I \otimes H_b + q \sum_{j=1}^N B_j \otimes \phi(g_j)$$

acting on the Hilbert space  $\mathcal{H} \otimes \mathcal{F}_b(\mathcal{K})$ , where  $q \in \mathbb{R}$  is a coupling parameter between the Bose field and the “particle”.

**Example 8.1** As is easily seen, the model  $H_{\text{NR}}$  in non-relativistic QED discussed in Section 2 is a concrete example of the GSB model with the following realizations (the unrenormalized version<sup>6</sup>):

$$\begin{aligned} N &= d, \\ \mathcal{H} &= L^2(\mathbb{R}^d), \quad \mathcal{K} = \oplus^{d-1} L^2(\mathbb{R}^d), \\ A &= H(m_0), \quad B_j = -\frac{p_j}{m_0 c}, \quad T = \hbar c \omega, \quad g_j = \frac{\sqrt{\hbar c} \hat{p} e_j}{\sqrt{\omega}}. \end{aligned}$$

As in the case of  $H_{\text{NR}}$ , we can consider a scaling limit for  $H_{\text{GSB}}$ . For that purpose, we consider the case where the particle Hamiltonian  $A$  is written

$$A = A_0 + A_1$$

with a non-negative self-adjoint operator  $A_0$  and a symmetric operator  $A_1$  on  $\mathcal{H}$ . We assume the following:

- (A.1)  $D(A_0) \subset D(A_1)$  and  $\|A_1 \psi\| \leq a \|A_0 \psi\| + b \|\psi\|$ ,  $\forall \psi \in D(A_0)$  for some constants  $a < 1$  and  $b \geq 0$ .
- (A.2) The set  $\{B_j\}_{j=1}^n$  is a set of strongly commuting self-adjoint operators on  $\mathcal{H}$ .
- (A.3) The self-adjoint operator  $A_0$  strongly commutes with  $B_j$ ,  $j = 1, \dots, N$ .
- (A.4)  $g_j \in D(T^{-1}) \cap D(T^{-1/2})$ ,  $j = 1, \dots, N$ , and  $\{g_j\}_{j=1}^N$  is linearly independent. Moreover, for all  $j, k = 1, \dots, N$ ,

$$\Lambda_{jk} := \langle T^{-1} g_j, T^{-1} g_k \rangle$$

and  $\langle T^{-1/2} g_j, T^{-1/2} g_k \rangle$  are real numbers.

Assumption (A.1) and the Kato-Rellich theorem[18, Theorem X.12] imply that  $A$  is self-adjoint and bounded below.

Condition (A.4) implies that the  $N \times N$  matrix

$$\Lambda := (\Lambda_{jk})$$

is strictly positive.

By condition (A.2), for all  $\mathbf{t} := (t_1, \dots, t_n) \in \mathbb{R}^n$ , the operator

$$B(\mathbf{t}) := \sum_{j=1}^n t_j B_j$$

---

<sup>6</sup>Also the renormalized version  $H_{\text{NR}}(\kappa)$  is included.

is essentially self-adjoint[5, Theorem 2.18], so that its closure  $\overline{B(\mathbf{t})}$  is self-adjoint and  $e^{\pm iq\overline{B(\mathbf{t})}}$  are unitary for all  $q \in \mathbb{R}$ . It is not so difficult to show that, for all  $\psi \in D(A_0)$  and all  $\mathbf{t}$ ,  $e^{iq\overline{B(\mathbf{t})}}\psi$  is in  $D(A_0)$  and  $e^{-iq\overline{B(\mathbf{t})}}A_1e^{iq\overline{B(\mathbf{t})}}\psi$  is strongly continuous in  $\mathbf{t}$  with

$$\|e^{-iq\overline{B(\mathbf{t})}}A_1e^{iq\overline{B(\mathbf{t})}}\psi\| \leq a\|A_0\psi\| + b\|\psi\|,$$

Here we have used (A.1) and (A.3) which implies that  $A_0e^{iq\overline{B(\mathbf{t})}} = e^{iq\overline{B(\mathbf{t})}}A_0$ . Hence we can define an operator  $A_1(q)$  with  $D(A_1(q)) = D(A_0)$  by

$$A_1(q)\psi := \frac{1}{\pi^{N/2}\sqrt{\det \Lambda}} \int_{\mathbb{R}^N} e^{-\langle \mathbf{t}, \Lambda^{-1}\mathbf{t} \rangle} e^{-iq\overline{B(\mathbf{t})}} A_1 e^{iq\overline{B(\mathbf{t})}} \psi d\mathbf{t}, \quad (8.1)$$

where the integral is taken in the sense of strong Riemann integral. It follows that  $A_1(q)$  is symmetric.

By condition (A.2), the operator

$$R_B := \frac{q^2}{2} \sum_{j,k=1}^N \langle T^{-1/2}g_j, T^{-1/2}g_k \rangle B_j B_k$$

is essentially self-adjoint[5, Theorem 2.18] and non-negative.

For each  $\kappa > 0$ , we define a renormalized version of  $H_{\text{GSB}}$  by

$$H_{\text{GSB}}(\kappa) := A \otimes I + \kappa R_B \otimes I + \kappa I \otimes H_b + \kappa q \sum_{j=1}^N B_j \otimes \phi(g_j).$$

By condition(A.4), the family  $\{\overline{\phi(iT^{-1}g_j)}\}_{j=1}^N$  of self-adjoint operators is strongly commuting[5, Theorem 4.43]. It follows from this fact and (A.2) that  $\{B_j \otimes \overline{\phi(iT^{-1}g_j)}\}_{j=1}^N$  is a set of strongly commuting self-adjoint operators. Thus the operator

$$L := \sum_{j=1}^N B_j \otimes \overline{\phi(iT^{-1}g_j)}$$

is essentially self-adjoint (see [9, Lemma 3.4]).

**Lemma 8.2** *For all  $\kappa > 0$ ,  $H_{\text{GSB}}(\kappa)$  is self-adjoint, bounded below and the operator equality*

$$e^{-iq\bar{L}}H_{\text{GSB}}(\kappa)e^{iq\bar{L}} = A_0 \otimes I + \kappa I \otimes H_b + H_1(q) \quad (8.2)$$

holds, where

$$H_1(q) := e^{-iq\bar{L}}(A_1 \otimes I)e^{iq\bar{L}}.$$

*Proof.* (Outline) One first proves (8.2) as an operator equality. We note that  $A_0 \otimes I$  strongly commutes with each  $B_j \otimes \overline{\phi(iT^{-1}g_j)}$  and hence with  $\bar{L}$ . Therefore  $e^{iq\bar{L}}$  leaves  $D(A_0 \otimes I)$  invariant, implying  $D(A_0 \otimes I) \subset D(H_1(q))$ . For all  $\Psi \in D(A_0 \otimes I)$ , we have

$$\|H_1(q)\Psi\| = \|A_1 \otimes I e^{iq\bar{L}}\Psi\| \leq a\|A_0 \otimes I\Psi\| + b\|\Psi\|.$$

Hence, for all  $\Psi \in D(A_0 \otimes I) \cap D(I \otimes H_b)$ ,

$$\|H_1(q)\Psi\| \leq a\|(A_0 \otimes I + \kappa I \otimes H_b)\Psi\| + b\|\Psi\|.$$

Since  $a < 1$ , it follows from the Kato-Rellich theorem that  $A_0 \otimes I + \kappa I \otimes H_b + H_1(q)$  is self-adjoint and bounded below. Thus the same holds for  $H_{\text{GSB}}(\kappa)$ .  $\blacksquare$

**Remark 8.3** In the previous paper [8] where  $\mathcal{K} = L^2(\mathbb{R}^\nu)$  and  $T$  is a multiplication operator on  $L^2(\mathbb{R}^\nu)$ , the unrenormalized Hamiltonian of the GSB model

$$H'_{\text{GSB}} := A \otimes I + I \otimes H_b + q \sum_{j=1}^N B_j \otimes \phi(g_j)$$

was considered with the following assumptions: (i) The particle Hamiltonian  $A$  is self-adjoint and bounded below. We denote by  $\mu_0$  the infimum of the spectrum of  $A$  and  $\tilde{A} := A_0 - \mu_0$  so that  $\tilde{A}$  is non-negative. (ii) Each  $B_j$  is  $\tilde{A}^{1/2}$ -bounded. (iii)  $g_j \in D(T^{-1/2})$  ( $j = 1, \dots, N$ ). Then it is shown that there exists a constant  $r_0 > 0$  such that, for all  $|q| < r_0$ ,  $H'_{\text{GSB}}$  is self-adjoint and bounded below. But, for the (renormalized) scaled Hamiltonian  $H_{\text{GSB}}(\kappa)$ , this method is not useful, because we have to consider the situation where  $\kappa \rightarrow \infty$  (then  $\kappa|q| \rightarrow \infty$ ) and hence we need conditions which ensure the self-adjointness of  $H_{\text{GSB}}(\kappa)$  for all sufficiently large  $\kappa$ . This is one of the reasons which led us to assumptions (A.1)–(A.4).

**Lemma 8.4** *Let*

$$A_{\text{eff}}(q) := A_0 + A_1(q).$$

*Then  $A_{\text{eff}}$  is self-adjoint with  $D(A_{\text{eff}}(q)) = D(A_0)$  and bounded below.*

*Proof.* We can show that  $\|A_1(q)\psi\| \leq a\|A_0\psi\| + b\|\psi\|$ ,  $\forall \psi \in D(A_0)$ , since

$$\frac{1}{\pi^{N/2} \sqrt{\det \Lambda}} \int_{\mathbb{R}^N} e^{-\langle \mathbf{t}, \Lambda^{-1} \mathbf{t} \rangle_{\mathbb{R}^N}} d\mathbf{t} = 1.$$

Hence, by Kato-Rellich theorem, the desired result follows.  $\blacksquare$

By Lemma 8.2 and an application of [4, Theorem 2.12], we can prove the following result on the scaling limit  $\kappa \rightarrow \infty$  for  $H_{\text{GSB}}(\kappa)$ :

**Theorem 8.5** *For all  $z \in \mathbb{C} \setminus \mathbb{R}$ ,*

$$s\text{-}\lim_{\kappa \rightarrow \infty} (H_{\text{GSB}}(\kappa) - z)^{-1} = e^{iq\bar{L}} (A_{\text{eff}}(q) - z)^{-1} \otimes P e^{-iq\bar{L}},$$

*where  $P$  is the orthogonal projection onto the one-dimensional subspace  $\{\alpha\Omega \mid \alpha \in \mathbb{C}\}$  with the Fock vacuum  $\Omega := \{1, 0, 0, \dots\} \in \mathcal{F}_b(\mathcal{K})$ .*

Theorem 8.5 shows that  $A_{\text{eff}}(q)$  can be regarded as an effective operator of  $H_{\text{GSB}}(\kappa)$  with  $A_1(q)$  being an effective “potential” for  $A_1$ .

**Example 8.6** Consider the case of Example 8.1. In this case, we take

$$A_0 = -\frac{\hbar^2}{2m(1)}\Delta, \quad A_1 = V,$$

where the potential  $V$  is assumed to satisfy (V.1) and (V.2) in Theorem 2.2. We also assume that  $\hat{\rho}$  is spherically symmetric. Then we have the following:

$$\Lambda_{jk} = \delta_{jk}C, \quad C := \frac{1}{\hbar c} \frac{d-1}{d} \left\| \frac{\hat{\rho}}{\omega^{3/2}} \right\|_{L^2(\mathbb{R}^d)}^2,$$

$$(A_1(q)\psi)(\mathbf{x}) = \frac{1}{\pi^{d/2}C^{d/2}} \int_{\mathbb{R}^d} e^{-\mathbf{t}^2/C} V \left( \mathbf{x} + \frac{q\hbar}{mc} \mathbf{t} \right) \psi(\mathbf{x}) d\mathbf{t}, \quad \psi \in D(\Delta).$$

By change of variables, one easily sees that

$$A_1(q)\psi = V_{\text{eff}}\psi, \quad \psi \in D(\Delta)$$

with  $V_{\text{eff}}$  given by (2.3). Hence  $\overline{A_1(q)} = V_{\text{eff}}$ .

Thus the operator  $A_{\text{eff}}(q)$  may be regarded as a perturbation of  $A$  which is physically meaningful to applications to quantum mechanics. To analyze spectral properties of  $A_{\text{eff}}(q)$  in comparison with those of  $A$  would be interesting.

**Remark 8.7** By condition (A.3), we have

$$A_{\text{eff}}(q) = \frac{1}{\pi^{N/2} \sqrt{\det \Lambda}} \int_{\mathbb{R}^N} e^{-\langle \mathbf{t}, \Lambda^{-1} \mathbf{t} \rangle_{\mathbb{R}^N}} e^{-iq\overline{B(\mathbf{t})}} A e^{iq\overline{B(\mathbf{t})}} d\mathbf{t}.$$

This suggests that a more general “effective” operator of  $A$  may be of the following form:

$$A_f(q) := \int_{\mathbb{R}^N} f(\mathbf{t}) e^{-iqC(\mathbf{t})} A e^{iqC(\mathbf{t})} d\mathbf{t},$$

where  $f$  is a real-valued Borel measurable function on  $\mathbb{R}^N$  such that  $\int_{\mathbb{R}^N} f(\mathbf{t}) d\mathbf{t} = 1$ ,  $C(\mathbf{t})$  is a self-adjoint operator on  $\mathcal{H}$ , and it is assumed that the right hand side is defined as a strong integral on a suitable dense subspace of  $\mathcal{H}$ . We propose to analyze properties of the operator  $A_f(q)$  including its self-adjointness.

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