The Yukawa potential: from Quantum Mechanics to Quantum Field Theory

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Abstract: Quantum Mechanics and Quantum Field Theory treatment for the scattering of two fermions is compared to find out an instantaneous effective Yukawa potential describing the interaction in the non relativistic limit. The comparison is made between the S matrix element for the scattering of 2 particles in time independent QM and the non relativistic limit of the scattering amplitude of 2 fermions due to a Yukawa interaction in QFT.

I. INTRODUCTION

The understanding of interactions between particles is essential for any theory that attempts to explain the fundamental structure and behaviour of matter, because it's this understanding which allows to make predictions that can be then observed (or not) in real life, and thus test a theory.
Quantum Mechanics explains that particles are sources of potential, and this potential covers some region of space that if another particle were to enter it, an atractive or repulsive interaction depending on the potential would then occur.
Quantum Field Theory describes interactions from the perspective of virtual particle exchange between the particles that are interacting, so instead of an instantaneous potential, there's a deeper understanding on how the propagation of information is carried.
The purpose of this work is to connect these two descriptions in the scattering of 2 fermions, in order to derive in the non relativistic limit one of the first discovered effective potentials between nucleons, the Yukawa potential.
This paper is written in natural units, $\hbar = c = 1$, and bold letters denote 3-vectors.

II. S MATRIX IN QM

Given a scattering process in QM due to some potential $\hat{V}$, we will work with a hamiltonian $\hat{H}$ that can be written as $\hat{H} = \hat{H}_0 + \hat{V}$, with $\hat{H}_0$ being the free (non interacting) hamiltonian, and the potential $\hat{V}$ causing the interaction. In this description, we want the solutions for the Schrödinger equation of $\hat{H}_0$ and $\hat{H}$ to have the same eigenvalue, meaning that when we connect the interaction $\hat{V}$ the energy stays the same as in the free case:

$$\hat{H}_0|\phi_a\rangle = E_a|\phi_a\rangle \tag{1}$$

$$\hat{H}|\psi_a\rangle = E_a|\psi_a\rangle \tag{2}$$

The scattering process will be described by how some initial $|\psi_i\rangle$ and final $|\psi_f\rangle$ interacting states with their respective free solutions $|\phi_i\rangle$ and $|\phi_f\rangle$ relate to each other. This connection will be made introducing an S matrix element as follows [1]:

$$S_{fi} = \langle\psi_f|\psi_i\rangle \tag{3}$$

Manipulating eqs. (1) and (2) in order to work out eq. (3), we can obtain the following relations between $|\psi_a\rangle$ and $|\phi_a\rangle$:

$$|\psi_a^{\pm}\rangle = |\phi_a\rangle + \frac{1}{(E_a - \hat{H}_0 \pm i\epsilon)}\hat{V}|\psi_a^{\pm}\rangle \tag{4}$$

$$|\psi_a^{\pm}\rangle = |\phi_a\rangle + \frac{1}{(E_a - \hat{H} \pm i\epsilon)}\hat{V}|\phi_a\rangle \tag{5}$$

Where the $\pm i\epsilon$ have been introduced to avoid the singularities provided that we take the limit $\epsilon \to 0$ at the end of the calculation. The initial (final) states are then defined as the ones with the $+$ $(-)$ sign for the physical reasons discussed in [1], [2], and so (3) now reads:

$$S_{fi} = \langle\psi_f^-|\psi_i^+\rangle \tag{6}$$

If we develop (6) by using the hermitian conjugate of eq. (5) for the $\langle\psi_f^-|$, and then using (4) on the first term for the $|\psi_i^+\rangle$, one gets:

$$\begin{aligned}\langle\psi_f^-|\psi_i^+\rangle &= \langle\phi_f|\phi_i\rangle - \frac{2i\epsilon}{(E_f - E_i)^2 + \epsilon^2}\langle\phi_f|\hat{V}|\psi_i^+\rangle \\ &= \delta_{fi} - 2\pi i\delta(E_f - E_i)\langle\phi_f|\hat{V}|\psi_i^+\rangle \end{aligned} \tag{7}$$

Where we have used the identity $\lim_{\epsilon\to\infty} \frac{1}{\pi}\frac{\epsilon}{x^2+\epsilon^2} = \delta(x)$. In eq. (7) we can substitute (4) as many times as we want, getting every time higher powers of $\hat{V}$ and thus an infinite series for the perturbative expansion of the scattering amplitudes.

III. 2 PARTICLE SCATTTERING IN QM

To first order in V, the second term in (7) becomes:

$$-2\pi i\delta(E_f - E_i)\langle\phi_f|\hat{V}|\phi_i\rangle \tag{8}$$

In orther to compare this result with the scattering of two fermions in QFT, we develop the term $\langle\phi_f|\hat{V}|\phi_i\rangle$ in (8) with a two particle state of plane waves, $|\phi_i\rangle = |\boldsymbol{p}_1, \boldsymbol{p}_2\rangle$, $|\phi_f\rangle = |\boldsymbol{p}'_1, \boldsymbol{p}'_2\rangle$ and a traslational invariant potential:

$$\int d^3x_1 \int d^3x_2 V(\boldsymbol{x}_1 - \boldsymbol{x}_2) e^{i(\boldsymbol{p}_1-\boldsymbol{p}'_1)\boldsymbol{x}_1} e^{i(\boldsymbol{p}_2-\boldsymbol{p}'_2)\boldsymbol{x}_2} \tag{9}$$

If the masses of the two particles are m_1, m_2, we can change $\{\boldsymbol{x}_1, \boldsymbol{x}_2\}$ for the center of mass and the relative coordinates $\{\boldsymbol{x}_{CM}, \boldsymbol{x}_R\}$:

$$\begin{aligned}\boldsymbol{x}_{CM} &= \frac{m_1\boldsymbol{x}_1 + m_2\boldsymbol{x}_2}{m_1 + m_2} \\ \boldsymbol{x}_R &= \boldsymbol{x}_1 - \boldsymbol{x}_2\end{aligned} \tag{10}$$

The determinant of the jacobian matrix for this transformation is equal to 1, and so (9) now reads:

$$\begin{aligned}&\int d^3x_{CM} e^{i(\boldsymbol{p}_1+\boldsymbol{p}_2-\boldsymbol{p}'_1-\boldsymbol{p}'_2)\boldsymbol{x}_{CM}} \\ &\times \int d^3x_R V(\boldsymbol{x}_R) e^{i\left(\frac{m_2(\boldsymbol{p}_1-\boldsymbol{p}'_1)}{m_1+m_2} - \frac{m_1(\boldsymbol{p}_2-\boldsymbol{p}'_2)}{m_1+m_2}\right)\boldsymbol{x}_R}\end{aligned} \tag{11}$$

From this change of variables we can then see that momentum conservation delta arises naturally, and defining $\boldsymbol{q} = \boldsymbol{p}_1 - \boldsymbol{p}'_1$ as the momentum transfered from one particle to the other, the integral over $\boldsymbol{x}_R$ becomes a Fourier transform of the potential. With this, the S matrix element to first order in $\hat{V}$ for the scattering of 2 particles is:

$$S_{fi} = \delta_{fi} - i(2\pi)^4\delta(E_f - E_i)\delta(\boldsymbol{p}_1 + \boldsymbol{p}_2 - \boldsymbol{p}'_1 - \boldsymbol{p}'_2)\widetilde{V}(\boldsymbol{q}) \tag{12}$$

IV. S MATRIX IN QFT

In QFT,we treat scattering from a time evolution perspective, meaning that the scattering amplitudes are calculated with the matrix element between the final state and the time evolution of the initial one.
Given a certain interacting hamiltonian $\hat{H} = \hat{H}_0 + \hat{H}_{int}$, where $\hat{H}_0$ describes the particle fields (Dirac, Klein Gordon, photons etc.) and $\hat{H}_{int}$ the interaction between them, when dealing with scattering it is useful to use the interaction picture. In this picture, states and operators are defined as:

$$|\psi(t)\rangle_I = e^{iH_0t}|\psi(t)\rangle_S \tag{13}$$
$$\hat{O}_I = e^{iH_0t}\hat{O}_S e^{-iH_0t} \tag{14}$$

Where the S subindex refers to the operator in the Schrödinger picture. Staying in the interaction picture, states evolve according to the following equation:

$$i\frac{d|\psi(t)\rangle_I}{dt} = \hat{H}_I|\psi(t)\rangle_I \tag{15}$$

If one tries to write the solution of (15) as $|\psi(t)\rangle_I = U_I(t, t_0)|\psi(t_0)\rangle_I$ with $U(t_0, t_0) = 1$, that would be:

$$\begin{aligned}U_I(t, t_0) &= 1 + (-i)\int_{t_0}^{t} dt_1 H_I(t_1) \\ &+ (-i)^2 \int_{t_0}^{t} dt_1 \int_{t_0}^{t_1} dt_1 dt_2 H_I(t_1)H_I(t_2) + \dots\end{aligned} \tag{16}$$

Eq.(16) can be written in a compact way introducing the time ordering operator T, which puts earlier time defined operators to the right and later ones to the right:

$$U(t, t_0) = T\{e^{-i\int_{t_0}^{t}\hat{H}_I(t')dt'}\} \tag{17}$$

Where since T is linear, the time ordering of the exponential is defined as the taylor series with each term time ordered. The S matrix is then defined as [3]:

$$\hat{S} = \lim_{t_\pm \to \pm\infty} U(t_-, t_+) \tag{18}$$

In QFT the S matrix is often decomposed separating the non interacting term from the interacting as follows:

$$S = \mathbb{1} + i\mathbb{T} \tag{19}$$

When we calculate scattering amplitudes, the identity term doesn't contribute because we consider processes in which the initial state is different from the final one.

V. FERMION-FERMION SCATTERING IN QFT

We want to look at the fermion-fermion scattering due to a Yukawa interaction:

$$H_I = \int d^3x g\phi\bar{\psi}\psi \tag{20}$$

Where g is the coupling constant of the interacting theory, with $|g| < 1$, and the ϕ, ψ, $\bar{\psi}$ fields are defined in the appendix. We will proceed to do this using Wick's theorem for the time ordered operator [3]. The free hamiltonian will be the sum of real Klein Gordon and Dirac hamiltonians, and every operator or state we write is in the interaction picture, thus we will drop the subindex I. Our initial and final states, $|i\rangle$ and $|f\rangle$, will be two fermions with definite momenta and spin with the relativistic normalization:

$$|i\rangle \equiv |p_1, p_2\rangle = \sqrt{2E_1}\sqrt{2E_2} b^{r\dagger}_{p_1} b^{s\dagger}_{p_2} |0\rangle \tag{21}$$
$$\langle f| \equiv \langle p'_1, p'_2| = \langle 0| b^{s'}_{p'_2} b^{r'}_{p'_1} \sqrt{2E'_2}\sqrt{2E'_1} \tag{22}$$

The scattering amplitude for the process is:

$$\langle p'_1, p'_2| T\{e^{-i\int d^4x \phi\bar{\psi}\psi}\} |p_1, p_2\rangle \tag{23}$$

The zero order term of the exponential series in (23) is clearly contributing to the $\mathbb{1}$ in (19), and so it is not relevant for the scattering process. Now if we look at the first order, this has only one $\phi \sim a^\dagger + a$, and these creation and anihilation operators conmute with the $b's$ and $c's$ because they create and anihilate different particles. Since our initial and final states don't have any bosons in them, this would end up with either $a|0\rangle$ or $\langle 0|a^\dagger$, giving it zero contribution.

We see then, the leading contribution is the second order:

$$\frac{(-ig)^2}{2!} \int d^4x \int d^4y\, T\Big\{ \big(\phi\bar{\psi}\psi\big)_x \big(\phi\bar{\psi}\psi\big)_y \Big\} \tag{24}$$

The only contribution from the time ordered product to the physical proccess is the one that anihilates two fermions on the right and two fermions on the left [3], [4]. Plugging it into (23) and omitting the integrals over x and y, we get the following expression:

$$\langle p'_1, p'_2| (-ig)^2 : \big(\phi\bar{\psi}\psi\big)_x \big(\phi\bar{\psi}\psi\big)_y : |p_1, p_2\rangle \tag{25}$$

The scalar propagator (see (39)) that arises from the contracion of the ϕ's won't be shown in the expressions until it is needed. The momenta and spins for fields (KG particles not included) are written as k_1, k_2, k_3, k_4 and l_1, l_2, l_3, l_4 from left to right. Expression (25), omitting the integrals and the constant factors, once it's normal ordered, looks like:

$$e^{i(k_1-k_2)x} e^{i(k_3-k_4)y} \big[\bar{u}_{l_1}(k_1) u_{l_2}(k_2)\big] \big[\bar{u}_{l_3}(k_3) u_{l_4}(k_4)\big]$$
$$\times (-1) \langle 0| b^{s'}_{p'_2} b^{r'}_{p'_1} \Big(b^{l_1\dagger}_{k_1} b^{l_3\dagger}_{k_3} b^{l_2}_{k_2} b^{l_4}_{k_4} \Big) b^{r\dagger}_{p_1} b^{s\dagger}_{p_2} |0\rangle \tag{26}$$

The minus sign comes from conmuting $b^{l_2}_{k_2}$ and $b^{l_3\dagger}_{k_3}$ and the "[ ]" are to denote that the Dirac spinor indexes are contracted inside of it. The minus sign will be left out until the final result. Using the anticonmutation relations (40), we move the b and $b^\dagger$ operators from the fields of the interaction through the $b^\dagger$'s and b's of the initial and final state. After performing then the integrals over all the k's, the expression left is:

$$\big[\bar{u}_{r'}(p'_1) u_s(p_2)\big] \big[\bar{u}_{s'}(p'_2) u_r(p_1)\big] e^{i(p'_1-p_2)x + i(p'_2-p_1)y}$$
$$- \big[\bar{u}_{r'}(p'_1) u_r(p_1)\big] \big[\bar{u}_{s'}(p'_2) u_s(p_2)\big] e^{i(p'_2-p_2)x + i(p'_1-p_1)y}$$
$$+ (x \leftrightarrow y) \tag{27}$$

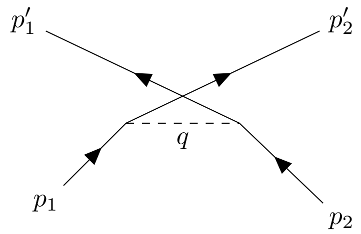
We now use the symetry of the propagator $D_F(x-y) = D_F(y-x)$ and multiply the first two terms by $D_F(y-x)$ and the last two by $D_F(x-y)$ so when we integrate x and y the contributions are the same for (x,y) as for (y,x). Plugging the exponentials from the propagator as mentioned and performing the itegrals over x and y, omitting the integral over the momentum q of the propagator, (27) becomes:

$$\big[\bar{u}_{r'}(p'_1) u_s(p_2)\big] \big[\bar{u}_{s'}(p'_2) u_r(p_1)\big] \times$$
$$(2\pi)^8 2\delta(p'_1 - p_2 + q)\delta(p'_2 - p_1 - q)$$
$$- \big[\bar{u}_{r'}(p'_1) u_r(p_1)\big] \big[\bar{u}_{s'}(p'_2) u_s(p_2)\big] \times$$
$$(2\pi)^8 2\delta(p'_2 - p_2 + q)\delta(p'_1 - p_1 - q) \tag{28}$$

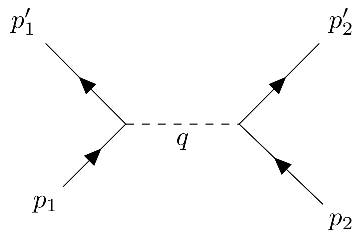
This 2 factor that appears from our convenient use of $D_F(x-y)$ cancels the 1/2! from (24). Performing the integral over q with either of the deltas that appear on each term, the final amplitude is:

$$ig^2 (2\pi)^4 \delta(p_1 + p_2 - p'_1 - p'_2) \times$$
$$\Bigg(\frac{[\bar{u}_{r'}(p'_1) u_s(p_2)][\bar{u}_{s'}(p'_2) u_r(p_1)]}{(p'_2 - p_1)^2 - m_\phi^2} - \frac{[\bar{u}_{r'}(p'_1) u_r(p_1)][\bar{u}_{s'}(p'_2) u_s(p_2)]}{(p_1 - p'_1)^2 - m_\phi^2} \Bigg) \tag{29}$$

This would correspond to the two possible diagrams of two identical fermions that interact exchanging a virtual scalar of mass m_ϕ. The diagram for the first term in (29) is:



And the diagram for the second one is:



VI. OBTENTION OF THE YUKAWA POTENTIAL

So far we have computed the amplitude for the scattering of 2 particles in QM, and the amplitude for the scattering of 2 fermions in QFT formalism. In orther to compare them, a few aspects must be taken into consideration:

Notice how the amplitude in QFT has 2 different terms, and in QM we only have one. In the QM case, when we fixed $|\boldsymbol{p}_1', \boldsymbol{p}_2'\rangle$ and $|\boldsymbol{p}_1, \boldsymbol{p}_2\rangle$ as the final and initial state, we forced the particles to be distinguishable, since momenta $\{\boldsymbol{p}_1', \boldsymbol{p}_1\}$ were assigned to the coordinate $\boldsymbol{x}_1$, and $\{\boldsymbol{p}_2', \boldsymbol{p}_2\}$ to $\boldsymbol{x}_2$. The calculation done in QFT doesn't have this restriction, as can be seen diagramatically.
If we had distinguishable fermions in QFT, as happens to be the case of QM, the only contribution to the amplitude would be the second term in (29), corresponding to the second diagram, which is the term that from now on we will consider for the comparison between the two amplitudes.

The amplitude in QFT must also be taken in the non relativistic limit, meaning that $m \gg \boldsymbol{p}$. We shall consider only terms to the lowest order in momenta, so $p^0 = m + \boldsymbol{p}^2/2m \to m$, and then for any 4-momentum $p = (m, \boldsymbol{p})$. In this limit, we have:

$$u_\lambda(p) \to \sqrt{m}\begin{pmatrix}\xi_\lambda \\ \xi_\lambda\end{pmatrix} \tag{30}$$

$$(p_1 - p_1')^2 = 0 - |\boldsymbol{p}_1 - \boldsymbol{p}_1'|^2 \equiv -\boldsymbol{q}^2 \tag{31}$$

Where ξ are two component spinors that satisfy $\xi_\lambda^\dagger \xi_{\lambda'} = \delta_{\lambda\lambda'}$, and $\boldsymbol{q}$ is the 3-momentum carried by the virtual scalar. The scattering amplitude then reads:

$$-ig^2(2\pi)^4\delta(p_1 + p_2 - p_1' - p_2')(2m)^2\left(\frac{\delta_{r'r}\delta_{s's}}{-\boldsymbol{q}^2 - m_\phi^2}\right) \tag{32}$$

Imposing spin conservation in this last expresion, we can finally connect it to the last term on the r.h.s. of (12):

$$\widetilde{V}(\boldsymbol{q}) = (2m)^2\frac{-g^2}{\boldsymbol{q}^2 + m_\phi^2} \tag{33}$$

The factor $(2m)^2$ comes from the relativistic normalization of our states in QFT: $\langle p|p'\rangle_{QFT} = 2E\delta(\boldsymbol{p} - \boldsymbol{p}')$, with $E = m$. We are now comparing it with an equation that has been computed with the non relativistic QM normalization $\langle p|p'\rangle = \delta(\boldsymbol{p} - \boldsymbol{p}')$, so the factor $(2m)^2$ should be dropped.
The integration on $\boldsymbol{q}$ can now be performed:

$$\begin{aligned} V(r) &= -g^2\int\frac{d^3q}{(2\pi)^3}\frac{e^{i\boldsymbol{qr}}}{\boldsymbol{q}^2 + m_\phi^2} \\ &= \frac{-g^2}{(2\pi)^2}\int_0^\infty d|\boldsymbol{q}|\frac{|\boldsymbol{q}|^2}{|\boldsymbol{q}|^2 + m_\phi^2}\int_{-1}^1 d(\cos\theta)\sin\theta e^{i|\boldsymbol{q}|r\cos\theta} \\ &= \frac{-g^2}{(2\pi)^2}\int_0^\infty d|\boldsymbol{q}|\frac{|\boldsymbol{q}|}{|\boldsymbol{q}|^2 + m_\phi^2}\frac{1}{ir}(e^{i|\boldsymbol{q}|r} - e^{-i|\boldsymbol{q}|r}) \\ &= \frac{ig^2}{(2\pi)^2 r}\int_{-\infty}^\infty d|\boldsymbol{q}|\frac{|\boldsymbol{q}|}{|\boldsymbol{q}|^2 + m_\phi^2}e^{i|\boldsymbol{q}|r} \end{aligned} \tag{34}$$

Closing the countour integral in the upper half of the complex plane:

$$V(r) = -\frac{g^2}{4\pi r}e^{-m_\phi r} \tag{35}$$

This is, the Yukawa potential arises as the lowest order effective potential for the interaction between two distinguible fermions in the non relativistic limit.

VII. CONCLUSIONS

We have seen how the scattering of two fermions due to a short range potential V is treated in time independent Quantum Mechanics, and calculated the S matrix element to first order in V. QFT formalism for interactions has been introduced, and we have obtained the scattering amplitude to leading order for 2 fermions that interact exchanging a Klein Gordon scalar.
We have seen that in the non relativistic limit and considering distinguible particles, these two results can be matched if the Yukawa potential is substituted as the potential causing the interaction in QM.

VIII. APPENDIX

A. Klein Gordon and Dirac fields

The Klein Gordon equation $(\partial_\mu\partial^\mu + m^2)\phi = 0$ describes a field for neutral bosons of spin 0. Its solution is:

$$\hat{\phi}(x) = \int\frac{d^3p}{(2\pi)^3}\frac{1}{\sqrt{2E}}\left(\hat{a}_p e^{-ipx} + \hat{a}_p^\dagger e^{ipx}\right) \tag{36}$$

Where $\hat{a}$ and $\hat{a}^\dagger$ are the anihilation and creation operators for Klein Gordon bosons, and all of the conmutators between them are zero except:

$$\left[\hat{a}_{\vec{p}}, \hat{a}_{\vec{q}}^\dagger\right] = (2\pi)^3\delta(\vec{p} - \vec{q}) \tag{37}$$

The Dirac equation $(i\partial\!\!\!/ - m)\psi = 0$ and its adjoint $(i\partial\!\!\!/ + m)\bar{\psi} = 0$ describe a field for spin 1/2 fermions, containing particles and antiparticles. The solutions for these two equations are:

$$\hat{\psi}(x) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E}} \sum_\lambda \left(e^{-ipx} u_\lambda(p) \hat{b}_p^\lambda + e^{ipx} v_\lambda(p) \hat{c}_p^{\lambda\dagger} \right) \quad (38)$$

$$\bar{\psi}(x) = \int \frac{d^3p}{(2\pi)^3} \frac{1}{\sqrt{2E}} \sum_\lambda \left(e^{ipx} \bar{u}_\lambda(p) \hat{b}_p^{\lambda\dagger} + e^{-ipx} \bar{v}_\lambda(p) \hat{c}_p^\lambda \right) \quad (39)$$

Where $\hat{b}$ ($\hat{c}$) and $\hat{b}^\dagger$ ($\hat{c}^\dagger$) are the anihilation and creation operators for Dirac particles (antiparticles). Since these are fermions, obey anticonmutation relations, and all of the anticonmutators between them are zero except:

$$\left\{ \hat{b}_p^\lambda, \hat{b}_q^{\lambda'\dagger} \right\} = \left\{ \hat{c}_p^\lambda, \hat{c}_q^{\lambda'\dagger} \right\} = (2\pi)^3 \delta(\vec{p} - \vec{q}) \delta^{\lambda\lambda'} \quad (40)$$

B. Propagator of the Klein Gordon field

The Wick contraction of two Klein Gordon fields at different times gives the propagator of the KG field (scalar). Written as a 4-momentum integral, this propagator is:

$$\phi_x \phi_y = D_F(x-y) = \int \frac{d^4p}{(2\pi)^4} \frac{ie^{-ip(x-y)}}{p^2 - m_\phi^2 + i\epsilon} \quad (41)$$

It is this exchange of a virtual ϕ particle between the two fermions expressed with this propagator which gives rise to the interaction and, in an effective way, to a non relativistic instantaneous Yukawa potential, as in (33).

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[1] L. de la Peña, *Introduction to Quantum Mechanics*, 3d ed. , Mexico: FCE, UNAM, 2006.

[2] M. Gell-Mann and L. Goldberger, "Formal Theory of Scattering". Phys. Rev. **91**, n°2 : 398–408 (1953).

[3] M. E. Peskin and D. V. Schroeder, *An introduction to Quantum Field Theory*, USA, Westview Press, 1995.

[4] D. Tong, Lecture notes on Quantum Field Theory, Cambridge, 2006-2007.

[5] J. J. Sakurai, *Modern quantum mechanics*, Addison-Wesley Pub. Co, 1994