



HAL
open science

WEAK INTERACTIONS IN A BACKGROUND OF A UNIFORM MAGNETIC FIELD. A MATHEMATICAL MODEL FOR THE INVERSE β DECAY.I.

Jean-Claude Guillot

► **To cite this version:**

Jean-Claude Guillot. WEAK INTERACTIONS IN A BACKGROUND OF A UNIFORM MAGNETIC FIELD. A MATHEMATICAL MODEL FOR THE INVERSE β DECAY.I.. 2017. hal-01585239

HAL Id: hal-01585239

<https://hal.science/hal-01585239v1>

Preprint submitted on 11 Sep 2017

HAL is a multi-disciplinary open access archive for the deposit and dissemination of scientific research documents, whether they are published or not. The documents may come from teaching and research institutions in France or abroad, or from public or private research centers.

L'archive ouverte pluridisciplinaire **HAL**, est destinée au dépôt et à la diffusion de documents scientifiques de niveau recherche, publiés ou non, émanant des établissements d'enseignement et de recherche français ou étrangers, des laboratoires publics ou privés.

A MATHEMATICAL MODEL FOR THE WEAK INTERACTIONS IN A BACKGROUND OF A UNIFORM MAGNETIC FIELD.I.THE INVERSE β DECAY.

JEAN-CLAUDE GUILLOT

In memory of Raymond Stora

ABSTRACT. In this paper we consider a mathematical model for the inverse β decay in a uniform magnetic field. With this model we associate a Hamiltonian with cutoffs in an appropriate Fock space. No infrared regularization is assumed. The Hamiltonian is self-adjoint and has a unique ground state. We study the essential spectrum and determine the spectrum. The coupling constant is supposed sufficiently small.

CONTENTS

1. Introduction.	2
2. The quantization of the Dirac fields for the electrons and the protons in a uniform magnetic field.	3
2.1. Eigenfunctions of the reduced Dirac operator for the electrons.	5
2.2. Eigenfunctions of the reduced Dirac operator for the protons.	7
2.3. Fock spaces for electrons, positrons, protons and antiprotons in a uniform magnetic field.	10
2.4. Quantized Dirac fields for the electrons and the protons in a uniform magnetic field.	11
3. The quantization of the Dirac fields for the neutrons and the neutrinos in helicity formalism.	13
3.1. The quantization of the Dirac field for the neutron in helicity formalism.	13
3.2. The quantization of the Dirac field for the neutrino.	16
4. The Hamiltonian of the Model.	19
4.1. The free Hamiltonian.	19
4.2. The Interaction.	22
5. Existence of a unique ground state for the Hamiltonian H.	29
5.1. The neutrino infrared cutoff Hamiltonians and the existence of a spectral gap	29
5.2. Proof of the existence of a ground state	32
5.3. Uniqueness of a ground state of the Hamiltonian H	34

2010 *Mathematics Subject Classification.* Primary 81Q10; Secondary 81V15, 81V10.

Key words and phrases. Weak Interactions, Uniform Magnetic Field, Beta decay Spectral Theory, Essential Spectrum, Unique Ground State.

1. INTRODUCTION.

A supernova is initiated by the collapse of a stellar core which leads to the formation of a protoneutron star which may be formed with strong magnetic fields typically of order 10^{16} Gauss. It turns out that the protoneutron star leads to the formation of a neutron star in a very short time during which almost all the gravitational binding energy of the protoneutron star is emitted in neutrinos and antineutrinos of each type. Neutron stars have strong magnetic fields of order 10^{12} Gauss. Thus neutrinos interactions are of great importance because of their capacity to serve as mediators for the transport and loss of energy and the following processes, the so-called "Urca" ones or inverse β decays in Physics,

$$(1.1) \quad \nu_e + n \rightleftharpoons e_- + p$$

$$(1.2) \quad \bar{\nu}_e + p \rightleftharpoons e_+ + n$$

play an essential role in those phenomena and they are associated with the β decay

$$(1.3) \quad n \rightarrow p + e_+ + \bar{\nu}_e$$

Here e_- (resp. e_+) is an electron (resp. a positron). p is a proton and n a neutron. ν_e and $\bar{\nu}_e$ are the neutrino and the antineutrino associated with the electron.

See [10], [11], [14], [17] and references therein.

We only consider here high-energy neutrinos and antineutrinos which are indeed ultrarelativistic particles whose mass is zero or in anyway negligible.

Due to the large magnetic field strengths involved, it is quite fundamental to study the processes (1.1) and (1.2) in the presence of magnetic fields.

These realistic fields may be very complicated in their structure but we assume these fields to be locally uniform which is a very good hypothesis because the range of the weak interactions is very short. Our aim is to study the processes (1.1) and (1.2) in a background of a uniform magnetic field.

Throughout this work we restrict ourselves to the study of processes (1.1), the study of processes (1.2) and (1.3) would be quite similar.

The advantage of a uniform magnetic field is that, in presence of this field, Dirac equation can be exactly solved. We can then quantize the corresponding field by using the canonical formalism and use the Fermi's Hamiltonian for the β decay in order to study the processes (1.1).

Throughout this paper we choose the units such that $c = \hbar = 1$.

In this paper we consider a mathematical model for the process (1.1) in a uniform magnetic field based on a Fock space for electrons, protons, neutrons and neutrinos and on a Hamiltonian with cutoffs suggested by the

Fermi's Hamiltonian for the β decay. The Fock space will also involve the antiparticles of the electrons and of the protons. No infrared regularization is assumed. We neglect the anomalous magnetic moments of the particles. Due to the cutoffs the relativistic invariance is broken.

We study the essential spectrum of the Hamiltonian and we prove the existence of a unique ground state. The spectrum of the Hamiltonian is identical to its essential spectrum. Every result is obtained for a sufficiently small coupling constant. We adapt to our case the proofs given in [4] and [7].

In an another paper we shall study the scattering theory and the absolutely continuous spectrum of the Hamiltonian.

The paper is organized as follows. In the next two sections we quantize the Dirac fields for electrons, protons and their antiparticles in a uniform magnetic field. In the third section we quantize the Dirac fields for free neutrons, neutrinos and their antiparticles in helicity formalism. The self-adjoint Hamiltonian of the model is defined in the fourth section. We then study the essential spectrum and prove the existence of a unique ground state.

2. THE QUANTIZATION OF THE DIRAC FIELDS FOR THE ELECTRONS AND THE PROTONS IN A UNIFORM MAGNETIC FIELD.

In this paper we assume that the uniform classical background magnetic field in \mathbb{R}^3 is along the x^3 -direction of the coordinate axis. There are several choices of gauge vector potential giving rise to a magnetic field of magnitude $B > 0$ along the x^3 -direction. In this paper we choose the following vector potential $\mathbf{A}(\mathbf{x}) = (A^\mu(\mathbf{x}), \mu = 0, 1, 2, 3)$, where

$$(2.1) \quad A^0(\mathbf{x}) = A^2(\mathbf{x}) = A^3(\mathbf{x}) = 0, \quad A^1(\mathbf{x}) = -x^2 B$$

Here $\mathbf{x} = (x^1, x^2, x^3)$ in \mathbb{R}^3 .

We recall that we neglect the anomalous magnetic moments of the particles of spin $\frac{1}{2}$.

The Dirac equation for a particle of spin $\frac{1}{2}$ with mass $m > 0$ and charge e in a uniform magnetic field of magnitude $B > 0$ along the x^3 -direction with the choice of the gauge (2.1) and by neglecting its anomalous magnetic moment is given by

$$(2.2) \quad H_D(e) = \boldsymbol{\alpha} \cdot \left(\frac{1}{i} \nabla - e\mathbf{A} \right) + \beta m,$$

acting in the Hilbert space $L^2(\mathbb{R}^3, \mathbb{C}^4)$.

The scalar product in $L^2(\mathbb{R}^3, \mathbb{C}^4)$ is given by

$$(f, g) = \sum_{j=1}^4 \int_{\mathbb{R}^3} \overline{f(\mathbf{x})} g(\mathbf{x}) d^3 \mathbf{x}$$

We refer to [27] for a discussion of the Dirac operator.

Here $\alpha = (\alpha_1, \alpha_2, \alpha_3)$, β are the Dirac matrices in the standard form:

$$\beta = \begin{pmatrix} I & 0 \\ 0 & -I \end{pmatrix}, \quad \alpha_i = \begin{pmatrix} 0 & \sigma_i \\ \sigma_i & 0 \end{pmatrix}, \quad i = 1, 2, 3,$$

where σ_i are the usual Pauli matrices.

By [27, thm 4.3] $H_D(e)$ is essentially self-adjoint on $C_0^\infty(\mathbb{R}^3, \mathbb{C}^4)$. The spectrum of $H_D(e)$ is equal to

$$(2.3) \quad \text{spec}(H_D(e)) = (-\infty, -m] \cup [m, \infty)$$

The spectrum of $H_D(e)$ is absolutely continuous and its multiplicity is not uniform. There is a countable set of thresholds, denoted by S , where

$$(2.4) \quad S = (-s_n, s_n; n \in \mathbb{N})$$

with $s_n = \sqrt{m^2 + 2n|e|B}$. See [19].

We consider a spectral representation of $H_D(e)$ based on a complete set of generalized eigenfunctions of the continuous spectrum of $H_D(e)$. Those generalized eigenfunctions are well known. See [22]. In view of (2.1) we use the computation of the generalized eigenfunctions given by [21] and [9]. See also [17] and references therein.

Let (p^1, p^3) be the conjugate variables of (x^1, x^3) . By the Fourier transform in \mathbb{R}^2 we easily get

$$(2.5) \quad L^2(\mathbb{R}^3, \mathbb{C}^4) \simeq \int_{\mathbb{R}^2}^\oplus L^2(\mathbb{R}, \mathbb{C}^4) dp^1 dp^3.$$

and

$$(2.6) \quad H_D(e) \simeq \int_{\mathbb{R}^2}^\oplus H_D(e; p^1, p^3) dp^1 dp^3.$$

where

$$(2.7) \quad H_D(e; p^1, p^3) = \begin{pmatrix} m\sigma_0, & \sigma_1(p^1 - ex^2B) - i\sigma_2 \frac{d}{dx^2} + p^3\sigma_3 \\ \sigma_1(p^1 - ex^2B) - i\sigma_2 \frac{d}{dx^2} + p^3\sigma_3, & -m\sigma_0 \end{pmatrix}$$

Here σ_0 is the 2×2 unit matrix.

$H_D(e; p^1, p^3)$ is the reduced Dirac operator associated to $(e; p^1, p^3)$.

$H_D(e; p^1, p^3)$ is essentially self-adjoint on $C_0^\infty(\mathbb{R}, \mathbb{C}^4)$ and has a pure point spectrum which is symmetrical with respect to the origin.

Set

$$(2.8) \quad E_n(p^3)^2 = m^2 + (p^3)^2 + 2n|e|B, \quad n \geq 0$$

The positive spectrum of $H_D(e; p^1, p^3)$ is the set of eigenvalues $(E_n(p^3))_{n \geq 0}$ and the negative spectrum is the set of eigenvalues $(-E_n(p^3))_{n \geq 0}$. $E_0(p^3)$ and $-E_0(p^3)$ are simple eigenvalues and the multiplicity of $E_n(p^3)$ and $-E_n(p^3)$ is equal to 2 for $n \geq 1$.

Through out this work e will be the positive unit of charge taken to be equal to the proton charge.

We now give the eigenfunctions of $H_D(e; p^1, p^3)$ both for the electrons and for the protons. The eigenfunctions are labelled by $n \in \mathbb{N}$, $(p_1, p_2) \in \mathbb{R}^2$ and $s = \pm 1$. $n \in \mathbb{N}$ labels the n th Landau level. $s = \pm 1$ are the eigenvalues of σ_3 . The electrons and the protons in all Landau levels with $n \geq 1$ can have different spin polarizations $s = \pm 1$. However in the lowest Landau state $n = 0$ the electrons can only have the spin orientation given by $s = -1$ and the protons can only have the spin orientation given by $s = 1$.

2.1. Eigenfunctions of the reduced Dirac operator for the electrons.

We now compute the eigenfunctions of $H_D(-e; p^1, p^3)$ with $m = m_e$ where m_e is the mass of the electron.

$E_n^{(e)}(p^3)$ and $-E_n^{(e)}(p^3)$ will denote the eigenvalues of $H_D(-e; p^1, p^3)$ for the electrons. We have $E_n^{(e)}(p^3)^2 = m_e^2 + (p^3)^2 + 2neB$, $n \geq 0$.

2.1.1. Eigenfunctions of the electrons for positive eigenvalues.

For $n \geq 1$ $E_n^{(e)}(p^3)$ is of multiplicity two corresponding to $s = \pm 1$ and $E_0^{(e)}(p^3)$ is multiplicity one corresponding to $s = -1$.

Let $U_{\pm 1}^{(e)}(x^2, n, p^1, p^3)$ denote the eigenfunctions associated to $s = \pm 1$.

For $s = 1$ and $n \geq 1$ we have

$$(2.9) \quad U_{+1}^{(e)}(x^2, n, p^1, p^3) = \left(\frac{E_n^{(e)}(p^3) + m_e}{2E_n^{(e)}(p^3)} \right)^{\frac{1}{2}} \begin{pmatrix} I_{n-1}(\xi) \\ 0 \\ \frac{p^3}{E_n^{(e)}(p^3) + m_e} I_{n-1}(\xi) \\ -\frac{\sqrt{2neB}}{E_n^{(e)}(p^3) + m_e} I_n(\xi) \end{pmatrix}$$

where

$$(2.10) \quad \xi = \sqrt{eB} \left(x^2 - \frac{p^1}{eB} \right)$$

$$I_n(\xi) = \left(\frac{\sqrt{eB}}{n! 2^n \sqrt{\pi}} \right)^{\frac{1}{2}} \exp(-\xi^2/2) H_n(\xi).$$

Here $H_n(\xi)$ is the Hermite polynomial of order n and we define

$$(2.11) \quad I_{-1}(\xi) = 0$$

For $n = 0$ and $s = 1$ we set

$$U_{+1}^{(e)}(x^2, 0, p^1, p^3) = 0$$

For $s = -1$ and $n \geq 0$ we have

$$(2.12) \quad U_{-1}^{(e)}(x^2, n, p^1, p^3) = \left(\frac{E_n^{(e)}(p^3) + m_e}{2E_n^{(e)}(p^3)} \right)^{\frac{1}{2}} \begin{pmatrix} 0 \\ I_n(\xi) \\ -\frac{\sqrt{2neB}}{E_n^{(e)}(p^3) + m_e} I_{n-1}(\xi) \\ -\frac{p^3}{E_n^{(e)}(p^3) + m_e} I_n(\xi) \end{pmatrix}$$

Note that

$$(2.13) \quad \int dx^2 U_s^{(e)}(x^2, n, p^1, p^3)^\dagger U_{s'}^{(e)}(x^2, n, p^1, p^3) = \delta_{ss'}$$

where \dagger is the adjoint in \mathbb{C}^4 .

2.1.2. *Eigenfunctions of the electrons for negative eigenvalues.*

For $n \geq 1$ $-E_n^{(e)}(p^3)$ is of multiplicity two corresponding to $s = \pm 1$ and $-E_0^{(e)}(p^3)$ is multiplicity one corresponding to $s = -1$.

Let $V_{\pm 1}^{(e)}(x^2, n, p^1, p^3)$ denote the eigenfunctions associated with the eigenvalue $-E_n^{(ap)}(p^3)$ and with $s = \pm 1$.

For $s = 1$ and $n \geq 1$ we have

$$(2.14) \quad V_{+1}^{(e)}(x^2, n, p^1, p^3) = \left(\frac{E_n^{(e)}(p^3) + m_e}{2E_n^{(e)}(p^3)} \right)^{\frac{1}{2}} \begin{pmatrix} -\frac{p^3}{E_n^{(e)}(p^3) + m_e} I_{n-1}(\xi) \\ \frac{\sqrt{2neB}}{E_n^{(e)}(p^3) + m_e} I_n(\xi) \\ I_{n-1}(\xi) \\ 0 \end{pmatrix}$$

and for $n = 0$ we set

$$V_{+1}^{(e)}(x^2, 0, p^1, p^3) = 0$$

For $s = -1$ and $n \geq 0$ we have

$$(2.15) \quad V_{-1}^{(e)}(x^2, n, p^1, p^3) = \left(\frac{E_n^{(e)}(p^3) + m_e}{2E_n^{(e)}(p^3)} \right)^{\frac{1}{2}} \begin{pmatrix} \frac{\sqrt{2neB}}{E_n^{(e)}(p^3) + m_e} I_{n-1}(\xi) \\ \frac{p^3}{E_n^{(e)}(p^3) + m_e} I_n(\xi) \\ 0 \\ I_n(\xi) \end{pmatrix}$$

Note that

$$(2.16) \quad \int dx^2 V_s^{(e)}(x^2, n, p^1, p^3)^\dagger V_{s'}^{(e)}(x^2, n, p^1, p^3) = \delta_{ss'}$$

where \dagger is the adjoint in \mathbb{C}^4 .

The sets $(U_{\pm 1}^{(e)}(\cdot, n, p^1, p^3))_{(n, p^1, p^3)}$ and $(V_{\pm 1}^{(e)}(\cdot, n, p^1, p^3))_{(n, p^1, p^3)}$ of vectors in $L^2(\mathbb{R}, \mathbb{C}^4)$ are a orthonormal basis of $L^2(\mathbb{R}, \mathbb{C}^4)$.

This yields for $\Psi(\mathbf{x})$ in $L^2(\mathbb{R}^3, \mathbb{C}^4)$

$$(2.17) \quad \Psi(\mathbf{x}) = \frac{1}{2\pi} \sum_{s=\pm 1} \text{L.i.m} \left(\sum_{n \geq 0} \int_{\mathbb{R}^2} dp^1 dp^3 e^{(p^1 x^1 + p^3 x^3)} \left(c_s^{(e)}(n, p^1, p^3) U_s^{(e)}(x^2, n, p^1, p^3) + d_s^{(e)}(n, p^1, p^3) V_s^{(e)}(x^2, n, p^1, p^3) \right) \right).$$

where $c_{+1}^{(e)}(0, p^1, p^3) = d_{+1}^{(e)}(0, p^1, p^3) = 0$

Let $\widehat{\Psi}(x^2; p^1, p^3)$ be the Fourier transform of $\Psi(\cdot)$ with respect to x^1 and x^3 :

$$\widehat{\Psi}(x^2; p^1, p^3) = \text{L.i.m} \frac{1}{2\pi} \int_{\mathbb{R}^2} e^{-(p^1 x^1 + p^3 x^3)} \Psi(x^1, x^2, x^3) dx^1 dx^3$$

We have

$$(2.18) \quad \begin{aligned} c_s^{(e)}(n, p^1, p^3) &= \int_{\mathbb{R}} U_s^{(e)}(x^2, n, p^1, p^3)^\dagger \widehat{\Psi}(x^2; p^1, p^3) dx^2 \\ d_s^{(e)}(n, p^1, p^3) &= \int_{\mathbb{R}} V_s^{(e)}(x^2, n, p^1, p^3)^\dagger \widehat{\Psi}(x^2; p^1, p^3) dx^2. \end{aligned}$$

The complex coefficients $c_s^{(e)}(n, p^1, p^3)$ and $d_s^{(e)}(n, p^1, p^3)$ satisfy

$$(2.19) \quad \begin{aligned} \|\Psi(\cdot)\|_{L^2(\mathbb{R}^3, \mathbb{C}^4)}^2 &= \\ \sum_{s=\pm 1} \sum_{n \geq 0} \int &\left(|c_s^{(e)}(n, p^1, p^3)|^2 + |d_s^{(e)}(n, p^1, p^3)|^2 \right) dp^1 dp^3 < \infty. \end{aligned}$$

2.2. Eigenfunctions of the reduced Dirac operator for the protons.

We now compute the eigenfunctions of $H_D(e; p^1, p^3)$ with $m = m_p$.

$E_n^{(p)}(p^3)$ and $-E_n^{(p)}(p^3)$ denote the eigenvalues of $H_D(e; p^1, p^3)$ for the proton. We have $E_n^{(p)}(p^3)^2 = m_p^2 + (p^3)^2 + 2neB$, $n \geq 0$.

2.2.1. Eigenfunctions of the proton for positive eigenvalues.

For $n \geq 1$ $E_n^{(p)}(p^3)$ is of multiplicity two corresponding to $s = \pm 1$ and $E_0^{(p)}(p^3)$ is multiplicity one corresponding to $s = 1$.

Let $U_{\pm 1}^{(p)}(x^2, n, p^1, p^3)$ denote the eigenfunctions associated with the eigenvalue $E_n^{(p)}(p^3)$ and with $s = \pm 1$.

For $s = 1$ and $n \geq 0$ we have

$$(2.20) \quad U_{+1}^{(p)}(x^2, n, p^1, p^3) = \left(\frac{E_n^{(p)}(p^3) + m_p}{2E_n^{(p)}(p^3)} \right)^{\frac{1}{2}} \begin{pmatrix} I_n(\tilde{\xi}) \\ 0 \\ \frac{p^3}{E_n^{(p)}(p^3) + m_p} I_n(\tilde{\xi}) \\ \frac{\sqrt{2neB}}{E_n^{(p)}(p^3) + m_p} I_{n-1}(\tilde{\xi}) \end{pmatrix}$$

where

$$(2.21) \quad \begin{aligned} \tilde{\xi} &= \sqrt{eB} \left(x^2 + \frac{p^1}{eB} \right) \\ I_{-1}(\tilde{\xi}) &= 0. \end{aligned}$$

For $s = -1$ and $n \geq 1$ we have

$$(2.22) \quad U_{-1}^{(p)}(x^2, n, p^1, p^3) = \left(\frac{E_n^{(p)}(p^3) + m_p}{2E_n^{(p)}(p^3)} \right)^{\frac{1}{2}} \begin{pmatrix} 0 \\ I_{n-1}(\tilde{\xi}) \\ \frac{\sqrt{2neB}}{E_n^{(p)}(p^3) + m_p} I_n(\tilde{\xi}) \\ -\frac{p^3}{E_n^{(p)}(p^3) + m_p} I_{n-1}(\tilde{\xi}) \end{pmatrix}$$

For $n = 0$ and $s = -1$ we set

$$U_{-1}^{(p)}(x^2, 0, p^1, p^3) = 0.$$

Note that

$$\int dx^2 U_s^{(p)}(x^2, n, p^1, p^3)^\dagger U_{s'}^{(p)}(x^2, n, p^1, p^3) = \delta_{ss'}$$

where \dagger is the adjoint in \mathbb{C}^4 .

2.2.2. *Eigenfunctions of the proton for negative eigenvalues.*

For $n \geq 1$ $-E_n^{(p)}(p^3)$ is of multiplicity two corresponding to $s = \pm 1$ and $-E_0^{(p)}(p^3)$ is multiplicity one corresponding to $s = 1$.

Let $V_{\pm 1}^{(p)}(x^2, n, p^1, p^3)$ denote the eigenfunctions associated with the eigenvalue $-E_n^{(p)}(p^3)$ and with $s = \pm 1$.

For $s = 1$ and $n \geq 0$ we have

$$(2.23) \quad V_{+1}^{(p)}(x^2, n, p^1, p^3) = \left(\frac{E_n^{(p)}(p^3) + m_p}{2E_n^{(p)}(p^3)} \right)^{\frac{1}{2}} \begin{pmatrix} -\frac{p^3}{E_n^{(p)}(p^3) + m_p} I_{n-1}(\tilde{\xi}) \\ -\frac{\sqrt{2neB}}{E_n^{(p)}(p^3) + m_p} I_{n-1}(\tilde{\xi}) \\ I_n(\tilde{\xi}) \\ 0 \end{pmatrix}$$

For $s = -1$ and $n \geq 1$ we have

$$(2.24) \quad V_{-1}^{(p)}(x^2, n, p^1, p^3) = \left(\frac{E_n^{(p)}(p^3) + m_p}{2E_n^{(p)}(p^3)} \right)^{\frac{1}{2}} \begin{pmatrix} -\frac{\sqrt{2neB}}{E_n^{(p)}(p^3) + m_p} I_n(\tilde{\xi}) \\ \frac{p^3}{E_n^{(p)}(p^3) + m_p} I_{n-1}(\tilde{\xi}) \\ 0 \\ I_{n-1}(\tilde{\xi}) \end{pmatrix}$$

and for $n = 0$ and $s = -1$ we set

$$V_{-1}^{(p)}(x^2, 0, p^1, p^3) = 0$$

Note that

$$(2.25) \quad \int dx^2 V_s^{(p)}(x^2, n, p^1, p^3)^\dagger V_{s'}^{(p)}(x^2, n, p^1, p^3) = \delta_{ss'}$$

where \dagger is the adjoint in \mathbb{C}^4 .

The sets $(U_{\pm 1}^{(p)}(\cdot, n, p^1, p^3))_{(n, p^1, p^3)}$ and $(V_{\pm 1}^{(p)}(\cdot, n, p^1, p^3))_{(n, p^1, p^3)}$ of vectors in $L^2(\mathbb{R}, \mathbb{C}^4)$ are an orthonormal basis of $L^2(\mathbb{R}, \mathbb{C}^4)$.

This yields for $\Psi(\mathbf{x})$ in $L^2(\mathbb{R}^3, \mathbb{C}^4)$

(2.26)

$$\Psi(\mathbf{x}) = \frac{1}{2\pi} \sum_{s=\pm 1} \text{L.i.m} \left(\sum_{n \geq 0} \int_{\mathbb{R}^2} dp^1 dp^3 e^{(p^1 x^1 + p^3 x^3)} \left(c_s^{(p)}(n, p^1, p^3) U_s^{(p)}(x^2, n, p^1, p^3) + d_s^{(p)}(n, p^1, p^3) V_s^{(p)}(x^2, n, p^1, p^3) \right) \right).$$

where $c_{-1}^{(p)}(0, p^1, p^3) = d_{-1}^{(p)}(0, p^1, p^3) = 0$

The complex coefficients $c_s^{(p)}(n, p^1, p^3)$ and $d_s^{(p)}(n, p^1, p^3)$ satisfy

(2.27)

$$\|\Psi(\cdot)\|_{L^2(\mathbb{R}^3, \mathbb{C}^4)}^2 = \sum_{s=\pm 1} \sum_{n \geq 0} \int \left(|c_s^{(p)}(n, p^1, p^3)|^2 + |d_s^{(p)}(n, p^1, p^3)|^2 \right) dp^1 dp^3 < \infty.$$

We have

$$(2.28) \quad \begin{aligned} c_s^{(p)}(n, p^1, p^3) &= \int_{\mathbb{R}} U_s^{(p)}(x^2, n, p^1, p^3)^\dagger \widehat{\Psi}(x^2; p^1, p^3) dx^2 \\ d_s^{(p)}(n, p^1, p^3) &= \int_{\mathbb{R}} V_s^{(p)}(x^2, n, p^1, p^3)^\dagger \widehat{\Psi}(x^2; p^1, p^3) dx^2. \end{aligned}$$

2.2.3. *Eigenfunctions of the positron for positive eigenvalues.*

The generalized eigenfunctions for the positron, denoted by $U_{\pm 1}^{(-e)}(x^2, n, p^1, p^3)$, are obtained from $U_{\pm 1}^{(p)}(x^2, n, p^1, p^3)$ by substituting the mass of the electron m_e for m_p . The associated eigenvalues are denoted by $E_n^{(-e)}(p^3)$ with $E_n^{(-e)}(p^3)^2 = m_e^2 + (p^3)^2 + 2neB$, $n \geq 0$.

2.2.4. *Eigenfunctions of the positron for negative eigenvalues.*

The generalized eigenfunctions for the positron, associated with the eigenvalues $-E_n^{(-e)}(p^3)$ and denoted by $V_{\pm 1}^{(-e)}(x^2, n, p^1, p^3)$, are obtained from $V_{\pm 1}^{(p)}(x^2, n, p^1, p^3)$ by substituting the mass of the electron m_e for m_p .

2.2.5. *Eigenfunctions of the antiproton for positive eigenvalues.*

The generalized eigenfunctions for the antiproton, denoted by $U_{\pm 1}^{(-p)}(x^2, n, p^1, p^3)$, are obtained from $U_{\pm 1}^{(e)}(x^2, n, p^1, p^3)$ by substituting the mass of the proton m_p for m_e . The associated eigenvalues are denoted by $E_n^{(-p)}(p^3)$ with $E_n^{(-p)}(p^3)^2 = m_p^2 + (p^3)^2 + 2neB$, $n \geq 0$.

2.2.6. Eigenfunctions of the antiproton for negative eigenvalues.

The generalized eigenfunctions for the antiproton, associated with the eigenvalues $-E_n^{(-p)}(p^3)$ and denoted by $V_{\pm 1}^{(-p)}(x^2, n, p^1, p^3)$, are obtained from $V_{\pm 1}^{(e)}(x^2, n, p^1, p^3)$ by substituting the mass of the proton m_p for m_e .

2.3. Fock spaces for electrons, positrons, protons and antiprotons in a uniform magnetic field.

It follows from section 2.1 and 2.2 that (s, n, p^1, p^3) are quantum variables for the electrons, the positrons, the protons and the antiprotons in a uniform magnetic field.

Let $\xi_1 = (s, n, p_e^1, p_e^3)$ be the quantum variables of a electron and of a positron and let $\xi_2 = (s, n, p_p^1, p_p^3)$ be the quantum variables of a proton and of an antiproton.

We set $\Gamma_1 = \{-1, 1\} \times \mathbb{N} \times \mathbb{R}^2$ for the configuration space for both the electrons, the positrons, the protons and the antiprotons. $L^2(\Gamma_1)$ is the Hilbert space associated to each species of fermions.

We have, by (2.17), (2.18), (2.19), (2.26), (2.27) and (2.28),

$$(2.29) \quad L^2(\Gamma_1) = l^2(L^2(\mathbb{R}^2)) \oplus l^2(L^2(\mathbb{R}^2))$$

Let $\mathfrak{F}_{(e)}$ and $\mathfrak{F}_{(-e)}$ denote the Fock spaces for the electrons and the positrons respectively and let $\mathfrak{F}_{(p)}$ and $\mathfrak{F}_{(-p)}$ denote the Fock spaces for the protons and the antiprotons respectively.

We have

$$(2.30) \quad \mathfrak{F}_{(e)} = \mathfrak{F}_{(-e)} = \mathfrak{F}_{(p)} = \mathfrak{F}_{(-p)} = \bigoplus_{n=0}^{\infty} \bigotimes_a^n L^2(\Gamma_1)$$

$\bigotimes_a^n L^2(\Gamma_1)$ is the antisymmetric n-th tensor power of $L^2(\Gamma_1)$.

$\Omega_{(\alpha)} = (1, 0, 0, 0, \dots)$ is the vacuum state in $\mathfrak{F}_{(\alpha)}$ for $\alpha = e, -e, p, -p$.

We shall use the notations

$$(2.31) \quad \int_{\Gamma_1} d\xi_1 = \sum_{s=\pm 1} \sum_{n \geq 0} \int_{\mathbb{R}^2} dp_e^1 dp_e^3$$

$$\int_{\Gamma_1} d\xi_2 = \sum_{s=\pm 1} \sum_{n \geq 0} \int_{\mathbb{R}^2} dp_p^1 dp_p^3.$$

Set $\epsilon = \pm$.

$b_\epsilon(\xi_j)$ (resp. $b_\epsilon^*(\xi_j)$) are the annihilation (resp. creation) operators for the electron when $j = 1$ and for the proton when $j = 2$ if $\epsilon = +$.

$b_\epsilon(\xi_j)$ (resp. $b_\epsilon^*(\xi_j)$) are the annihilation (resp. creation) operators for the positron when $j = 1$ and for the antiproton when $j = 2$ if $\epsilon = -$.

The operators $b_\epsilon(\xi_j)$ and $b_\epsilon^*(\xi_j)$ fulfil the usual anticommutation relations (CAR) (see [28]).

In addition, following the convention described in [28, Section 4.1] and [28, Section 4.2], we assume that the fermionic creation and annihilation

operators of different species of particles anticommute (see [6] arXiv for explicit definitions). In our case this property will be verified by the creation and annihilation operators for the electrons, the protons, the neutrons, the neutrinos and their respective antiparticles.

Therefore the following anticommutation relations hold for $j = 1, 2$

$$(2.32) \quad \begin{aligned} \{b_\epsilon(\xi_j), b_{\epsilon'}^*(\xi'_j)\} &= \delta_{\epsilon\epsilon'}\delta(\xi_j - \xi'_j) , \\ \{b_\epsilon^\sharp(\xi_1), b_{\epsilon'}^\sharp(\xi_2)\} &= 0. \end{aligned}$$

where $\{b, b'\} = bb' + b'b$ and $b^\sharp = b$ or b^* .

Recall that for $\varphi \in L^2(\Gamma_1)$, the operators

$$(2.33) \quad \begin{aligned} b_{j,\epsilon}(\varphi) &= \int_{\Gamma_1} b_\epsilon(\xi_j)\overline{\varphi(\xi_j)}d\xi_j. \\ b_{j,\epsilon}^*(\varphi) &= \int_{\Gamma_1} b_\epsilon^*(\xi_j)\varphi(\xi_j)d\xi_j. \end{aligned}$$

are bounded operators on $\mathfrak{F}_{(e)}$ and $\mathfrak{F}_{(-e)}$ for $j = 1$ and on $\mathfrak{F}_{(p)}$ and $\mathfrak{F}_{(-p)}$ for $j = 2$ respectively satisfying

$$(2.34) \quad \|b_{j,\epsilon}^\sharp(\varphi)\| = \|\varphi\|_{L^2} ,$$

2.4. Quantized Dirac fields for the electrons and the protons in a uniform magnetic field.

We now consider the canonical quantization of the two classical fields (2.17) and (2.26).

Recall that the charge conjugation operator C is given, for every $\Psi(\mathbf{x})$, by

$$(2.35) \quad C \begin{pmatrix} \Psi_1(\mathbf{x}) \\ \Psi_2(\mathbf{x}) \\ \Psi_3(\mathbf{x}) \\ \Psi_4(\mathbf{x}) \end{pmatrix} = \begin{pmatrix} -\Psi_4^*(\mathbf{x}) \\ \Psi_3^*(\mathbf{x}) \\ \Psi_2^*(\mathbf{x}) \\ -\Psi_1^*(\mathbf{x}) \end{pmatrix}$$

Here $*$ is the complex conjugation.

Let $\Psi(\cdot)$ be locally in the domain of $H_D(e)$. We have

$$(2.36) \quad H_D(-e)C\Psi = EC\Psi \quad \text{if} \quad H_D(e)\Psi = -E\Psi$$

By (2.35) and (2.36) we obtain

$$(2.37) \quad \begin{aligned} (CV_{+1}^{(e)})(x^2, n, p^1, p^3) &= U_{-1}^{(-e)}(x^2, n, -p^1, -p^3) \quad \text{for } n \geq 1. \\ (CV_{-1}^{(e)})(x^2, n, p^1, p^3) &= -U_{+1}^{(-e)}(x^2, n, -p^1, -p^3) \quad \text{for } n \geq 0. \\ (CV_{+1}^{(p)})(x^2, n, p^1, p^3) &= U_{-1}^{(-p)}(x^2, n, -p^1, -p^3) \quad \text{for } n \geq 0. \\ (CV_{-1}^{(p)})(x^2, n, p^1, p^3) &= -U_{+1}^{(-p)}(x^2, n, -p^1, -p^3) \quad \text{for } n \geq 1. \end{aligned}$$

By (2.37) we set

$$(2.38) \quad \begin{aligned} U^{(e)}(x^2, \xi_1) &= U_s^{(e)}(x^2, n, p_e^1, p_e^3) \quad \text{for } \xi_1 = (s, n, p_e^1, p_e^3), n \geq 0. \\ W^{(e)}(x^2, \xi_1) &= V_{-1}^{(e)}(x^2, n, -p_e^1, -p_e^3) \quad \text{for } \xi_1 = (1, n, p_e^1, p_e^3), n \geq 0. \\ W^{(e)}(x^2, \xi_1) &= V_{+1}^{(e)}(x^2, n, -p_e^1, -p_e^3) \quad \text{for } \xi_1 = (-1, n, p_e^1, p_e^3), n \geq 1. \\ W^{(e)}(x^2, \xi_1) &= 0 \quad \text{for } \xi_1 = (-1, 0, p_e^1, p_e^3). \end{aligned}$$

By using (2.37) and (2.38) the symmetric of charge canonical quantization of the classical field (2.17) gives the following formal operator associated with the electron and denoted by $\Psi_{(e)}(\mathbf{x})$:

$$(2.39) \quad \begin{aligned} \Psi_{(e)}(\mathbf{x}) &= \frac{1}{2\pi} \int d\xi_1 \left(e^{i(p_e^1 x^1 + p_e^3 x^3)} U^{(e)}(x^2, \xi_1) b_+(\xi_1) \right. \\ &\quad \left. + e^{-i(p_e^1 x^1 + p_e^3 x^3)} W^{(e)}(x^2, \xi_1) b_-^*(\xi_1) \right). \end{aligned}$$

For a rigorous approach of the quantization see [13].

We further note that

$$(2.40) \quad \{\Psi_{(e)}(\mathbf{x}), \Psi_{(e)}(\mathbf{x}')^\dagger\} = \delta(\mathbf{x}, \mathbf{x}')$$

See [9].

By (2.37) we now set

$$(2.41) \quad \begin{aligned} U^{(p)}(x^2, \xi_2) &= U_s^{(p)}(x^2, n, p_p^1, p_p^3) \quad \text{for } \xi_2 = (s, n, -p_p^1, -p_p^3), n \geq 0. \\ W^{(p)}(x^2, \xi_2) &= V_{+1}^{(p)}(x^2, n, -p_p^1, -p_p^3) \quad \text{for } \xi_2 = (-1, n, p_p^1, p_p^3), n \geq 0. \\ W^{(p)}(x^2, \xi_2) &= V_{-1}^{(p)}(x^2, n, -p_p^1, -p_p^3) \quad \text{for } \xi_2 = (1, n, p_p^1, p_p^3), n \geq 1. \\ W^{(p)}(x^2, \xi_2) &= 0 \quad \text{when } \xi_2 = (1, 0, p_p^1, p_p^3). \end{aligned}$$

By using (2.37) and (2.41) the symmetric of charge canonical quantization of the classical field (2.26) gives the following formal operator associated to the proton and denoted by $\Psi_{(p)}(\mathbf{x})$:

$$(2.42) \quad \begin{aligned} \Psi_{(p)}(\mathbf{x}) &= \frac{1}{2\pi} \int d\xi_2 \left(e^{i(p_p^1 x^1 + p_p^3 x^3)} U^{(p)}(x^2, \xi_2) b_+(\xi_2) \right. \\ &\quad \left. + e^{-i(p_p^1 x^1 + p_p^3 x^3)} W^{(p)}(x^2, \xi_2) b_-^*(\xi_2) \right). \end{aligned}$$

We further note that

$$(2.43) \quad \{\Psi_{(p)}(\mathbf{x}), \Psi_{(p)}(\mathbf{x}')^\dagger\} = \delta(\mathbf{x} - \mathbf{x}')$$

See [9].

3. THE QUANTIZATION OF THE DIRAC FIELDS FOR THE NEUTRONS AND THE NEUTRINOS IN HELICITY FORMALISM.

As stated in the introduction we neglect the magnetic moment of the neutrons. Therefore neutrons and neutrinos are purely neutral particles without any electromagnetic interaction. We suppose that the neutrinos and antineutrinos are massless as in the Standard Model.

The quantized Dirac fields for free massive and massless particles of spin $\frac{1}{2}$ are well-known.

In this work we use the helicity formalism, for free particles. See, for example, [27], [23] and [18].

The helicity formalism for particles is associated with a spectral representation of the set of commuting self adjoint operators (\mathbf{P}, H^3) . $\mathbf{P} = (P^1, P^2, P^3)$ are the generators of space-translations and H^3 is the helicity operator $\frac{1}{2} \frac{\mathbf{P} \cdot \boldsymbol{\Sigma}}{|\mathbf{P}|}$ where $|\mathbf{P}| = (\sqrt{\sum_{i=1}^3 (P^i)^2})$ and $\boldsymbol{\Sigma} = (\Sigma^1, \Sigma^2, \Sigma^3)$ with for $j = 1, 2, 3$

$$(3.1) \quad \Sigma^j = \begin{pmatrix} \sigma_j & 0 \\ 0 & \sigma_j \end{pmatrix}$$

3.1. The quantization of the Dirac field for the neutron in helicity formalism.

The Dirac equation for the neutron of mass m_n is given by

$$(3.2) \quad H_D = \boldsymbol{\alpha} \cdot \frac{1}{i} \nabla + \beta m_n,$$

acting in the Hilbert space $L^2(\mathbb{R}^3, \mathbb{C}^4)$.

It follows from the Fourier transform that

$$(3.3) \quad \begin{aligned} L^2(\mathbb{R}^3, \mathbb{C}^4) &\simeq \int_{\mathbb{R}^3}^{\oplus} \mathbb{C}^4 d^3 p. \\ H_D &\simeq \int_{\mathbb{R}^3}^{\oplus} H_D(\mathbf{p}) d^3 p. \end{aligned}$$

where

$$(3.4) \quad H_D(\mathbf{p}) = \begin{pmatrix} m_n \sigma_0 & \boldsymbol{\sigma} \cdot \mathbf{p} \\ \boldsymbol{\sigma} \cdot \mathbf{p} & -m_n \sigma_0 \end{pmatrix}$$

Here σ_0 is the 2×2 unit matrix, $\boldsymbol{\sigma} = (\sigma_1, \sigma_2, \sigma_3)$ and $\mathbf{p} = (p^1, p^2, p^3)$ with $\boldsymbol{\sigma} \cdot \mathbf{p} = \sum_{j=1}^3 \sigma_j p^j$.

$H_D(\mathbf{p})$ has two eigenvalues $E^{(n)}(\mathbf{p})$ and $-E^{(n)}(\mathbf{p})$ where $E^{(n)}(\mathbf{p}) = \sqrt{|\mathbf{p}|^2 + m_n^2}$.

The helicity, denoted by $H_3(\mathbf{p})$, is given by

$$(3.5) \quad H_3(\mathbf{p}) = \frac{1}{2} \begin{pmatrix} \frac{\boldsymbol{\sigma} \cdot \mathbf{p}}{|\mathbf{p}|} & 0 \\ 0 & \frac{\boldsymbol{\sigma} \cdot \mathbf{p}}{|\mathbf{p}|} \end{pmatrix}$$

$H_3(\mathbf{p})$ commutes with $H_D(\mathbf{p})$ and has two eigenvalues $\frac{1}{2}$ and $-\frac{1}{2}$.
Set (see [27, Appendix.1.F.] and [23]) for $|\mathbf{p}| \neq p^3$

$$(3.6) \quad h_+(\mathbf{p}) = \frac{1}{\sqrt{2|\mathbf{p}|(|\mathbf{p}| - p^3)}} \begin{pmatrix} p^1 - ip^2 \\ |\mathbf{p}| - p^3 \end{pmatrix}$$

and

$$(3.7) \quad h_-(\mathbf{p}) = \frac{1}{\sqrt{2|\mathbf{p}|(|\mathbf{p}| - p^3)}} \begin{pmatrix} p^3 - |\mathbf{p}| \\ p^1 + ip^2 \end{pmatrix}$$

For $|\mathbf{p}| = p^3$ we set

$$h_+(\mathbf{p}) = \begin{pmatrix} 1 \\ 0 \end{pmatrix}$$

and

$$h_-(\mathbf{p}) = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$$

We have $(\sigma \cdot \mathbf{p})h_{\pm}(\mathbf{p}) = \pm|\mathbf{p}|h_{\pm}(\mathbf{p})$.

Let

$$(3.8) \quad a_{\pm}(\mathbf{p}) = \frac{1}{\sqrt{2}} \left(1 \pm \frac{m_n}{E^{(n)}(\mathbf{p})} \right)^{\frac{1}{2}}$$

The two eigenfunctions of the eigenvalue $E^{(n)}(\mathbf{p})$ associated with helicities $\frac{1}{2}$ and $-\frac{1}{2}$ are denoted by $U^{(n)}(\mathbf{p}, \pm\frac{1}{2})$ and are given by

$$(3.9) \quad U^{(n)}(\mathbf{p}, \pm\frac{1}{2}) = \begin{pmatrix} a_+(\mathbf{p})h_{\pm}(\mathbf{p}) \\ \pm a_-(\mathbf{p})h_{\pm}(\mathbf{p}) \end{pmatrix}$$

We now turn to the eigenfunctions for the eigenvalue $-E^{(n)}(\mathbf{p})$.

The two eigenfunctions associated with the eigenvalue $-E^{(n)}(\mathbf{p})$ and with helicities $\frac{1}{2}$ and $-\frac{1}{2}$ are denoted by $V^{(n)}(\mathbf{p}, \pm\frac{1}{2})$ and are given by

$$(3.10) \quad V^{(n)}(\mathbf{p}, \pm\frac{1}{2}) = \begin{pmatrix} \mp a_-(\mathbf{p})h_{\pm}(\mathbf{p}) \\ a_+(\mathbf{p})h_{\pm}(\mathbf{p}) \end{pmatrix}$$

The four vectors $U^{(n)}(\mathbf{p}, \pm\frac{1}{2})$ and $V^{(n)}(\mathbf{p}, \pm\frac{1}{2})$ are an orthonormal basis of \mathbb{C}^4 .

$U^{(n)}(\mathbf{p}, \pm\frac{1}{2})e^{i(\mathbf{p} \cdot \mathbf{x})}$ and $V^{(n)}(\mathbf{p}, \pm\frac{1}{2})e^{i(\mathbf{p} \cdot \mathbf{x})}$ is a complete set of generalized eigenfunctions of (3.2) with positive and negative eigenvalues $\pm E^{(n)}(\mathbf{p})$.

This yields for $\Psi(\mathbf{x})$ in $L^2(\mathbb{R}^3, \mathbb{C}^4)$

$$(3.11) \quad \Psi(\mathbf{x}) = \left(\frac{1}{2\pi} \right)^{\frac{3}{2}} \sum_{\lambda=\pm\frac{1}{2}} \text{L.i.m.} \left(\int_{\mathbb{R}^3} d^3\mathbf{p} e^{i(\mathbf{p} \cdot \mathbf{x})} \left(U^{(n)}(\mathbf{p}, \lambda)a(\mathbf{p}, \lambda) + V^{(n)}(\mathbf{p}, \lambda)c(\mathbf{p}, \lambda) \right) \right).$$

with

$$(3.12) \quad \|\Psi(\cdot)\|_{L^2(\mathbb{R}^3, \mathbb{C}^4)}^2 = \sum_{\lambda=\pm\frac{1}{2}} \int_{\mathbb{R}^3} d^3p (|a(\mathbf{p}, \lambda)|^2 + |c(\mathbf{p}, \lambda)|^2) < \infty.$$

3.1.1. **Fock space for the neutrons.** We recall that the neutron is not its own antiparticle.

Let $\xi_3 = (\mathbf{p}, \lambda)$ be the quantum variables of a neutron and an antineutron where $\mathbf{p} \in \mathbb{R}^3$ is the momentum and $\lambda \in \{-\frac{1}{2}, \frac{1}{2}\}$ is the helicity. We set $\Gamma_2 = \mathbb{R}^3 \times \{-\frac{1}{2}, \frac{1}{2}\}$ for the configuration space of the neutron and the antineutron.

Let $\mathfrak{F}_{(n)}$ and $\mathfrak{F}_{(\bar{n})}$ denote the Fock spaces for the neutrons and the antineutrons respectively.

We have

$$(3.13) \quad \mathfrak{F}_{(n)} = \mathfrak{F}_{(\bar{n})} = \bigoplus_{n=0}^{\infty} \bigotimes_a^n L^2(\Gamma_2)$$

$\bigotimes_a^n L^2(\Gamma_2)$ is the antisymmetric n-th tensor power of $L^2(\Gamma_2)$.

$\Omega_{(\beta)} = (1, 0, 0, 0, \dots)$ is the vacuum state in $\mathfrak{F}_{(\beta)}$ for $\beta = n, \bar{n}$.

In the sequel we shall use the notations

$$(3.14) \quad \int_{\Gamma_2} d\xi_3 = \sum_{\lambda=\pm\frac{1}{2}} \int_{\mathbb{R}^3} d^3p$$

$b_\epsilon(\xi_3)$ (resp. $b_\epsilon^*(\xi_3)$) is the annihilation (resp. creation) operator for the neutron if $\epsilon = +$ and for the antineutron if $\epsilon = -$.

The operators $b_\epsilon(\xi_3)$ and $b_\epsilon^*(\xi_3)$ fulfil the usual anticommutation relations (CAR) and they anticommute with $b_\epsilon^\sharp(\xi_j)$ for $j = 1, 2$ according to the convention described in [28, Section 4.1]. See [6] arXiv for explicit definitions.

Therefore the following anticommutation relations hold for $j = 1, 2$

$$(3.15) \quad \begin{aligned} \{b_\epsilon(\xi_3), b_{\epsilon'}^*(\xi'_3)\} &= \delta_{\epsilon\epsilon'} \delta(\xi_3 - \xi'_3) , \\ \{b_\epsilon^\sharp(\xi_3), b_{\epsilon'}^\sharp(\xi_j)\} &= 0. \end{aligned}$$

Recall that for $\varphi \in L^2(\Gamma_2)$, the operators

$$(3.16) \quad \begin{aligned} b_{3,\epsilon}(\varphi) &= \int_{\Gamma_2} b_\epsilon(\xi_3) \overline{\varphi(\xi_3)} d\xi_3. \\ b_{3,\epsilon}^*(\varphi) &= \int_{\Gamma_2} b_\epsilon^*(\xi_3) \varphi(\xi_3) d\xi_3. \end{aligned}$$

are bounded operators on $\mathfrak{F}_{(n)}$ and $\mathfrak{F}_{(\bar{n})}$ satisfying

$$(3.17) \quad \|b_{3,\epsilon}^\sharp(\varphi)\| = \|\varphi\|_{L^2}.$$

3.1.2. Quantized Dirac Field for the neutron in helicity formalism.

By (2.35) we get

$$(3.18) \quad \begin{aligned} C(V^{(n)}(\mathbf{p}, \frac{1}{2})) &= \left(-\frac{p^1 + ip^2}{|p^1 + ip^2|}\right) U^{(n)}(-\mathbf{p}, \frac{1}{2}) \\ C(V^{(n)}(\mathbf{p}, -\frac{1}{2})) &= \left(-\frac{p^1 - ip^2}{|p^1 + ip^2|}\right) U^{(n)}(-\mathbf{p}, -\frac{1}{2}) \end{aligned}$$

Setting

$$(3.19) \quad \begin{aligned} U^{(n)}(\xi_3) &= U^{(n)}(\mathbf{p}, \lambda) \\ W^{(n)}(\xi_3) &= V^{(n)}(-\mathbf{p}, \lambda). \end{aligned}$$

and applying the canonical quantization we obtain the following quantized Dirac field for the neutron:

$$(3.20) \quad \Psi_{(n)}(\mathbf{x}) = \left(\frac{1}{2\pi}\right)^{\frac{3}{2}} \int d\xi_3 \left(e^{i(\mathbf{p}\cdot\mathbf{x})} U^{(n)}(\xi_3) b_+(\xi_3) + e^{-i(\mathbf{p}\cdot\mathbf{x})} W^{(n)}(\xi_3) b_-^*(\xi_3) \right).$$

3.2. The quantization of the Dirac field for the neutrino.

Throughout this work we suppose that the neutrinos we consider are those associated with the electrons.

The Dirac equation for the neutrino is given by

$$(3.21) \quad H_D = \boldsymbol{\alpha} \cdot \frac{1}{i} \nabla,$$

acting in the Hilbert space $L^2(\mathbb{R}^3, \mathbb{C}^4)$.

By (3.3) it follows from the Fourier transform that

$$(3.22) \quad H_D \simeq \int_{\mathbb{R}^3}^{\oplus} H_D(\mathbf{p}) d^3p.$$

where

$$(3.23) \quad H_D(\mathbf{p}) = \begin{pmatrix} 0, & \sigma \cdot \mathbf{p} \\ \sigma \cdot \mathbf{p}, & 0 \end{pmatrix}$$

$H_D(\mathbf{p})$ has two eigenvalues $E^{(\nu)}(\mathbf{p})$ and $-E^{(\nu)}(\mathbf{p})$ where $E^{(\nu)}(\mathbf{p}) = |\mathbf{p}|$.

The helicity given by

$$\frac{1}{2} \gamma^5 = \frac{1}{2} \begin{pmatrix} 0 & I \\ I & 0 \end{pmatrix}$$

commutes with $H_D(\mathbf{p})$ and has two eigenvalues $\frac{1}{2}$ and $-\frac{1}{2}$.

The two eigenfunctions of the eigenvalue $E^{(\nu)}(\mathbf{p})$ associated with helicities $\frac{1}{2}$ and $-\frac{1}{2}$ are denoted by $U^{(\nu)}(\mathbf{p}, \pm\frac{1}{2})$. The two eigenfunctions of the

eigenvalue $-E^{(\nu)}(\mathbf{p})$ associated with helicities $\frac{1}{2}$ and $-\frac{1}{2}$ are denoted by $V^{(\nu)}(\mathbf{p}, \pm\frac{1}{2})$. They are given by

$$(3.24) \quad \begin{aligned} U^{(\nu)}(\mathbf{p}, \pm\frac{1}{2}) &= \frac{1}{\sqrt{2}} \begin{pmatrix} h_{\pm}(\mathbf{p}) \\ \pm h_{\pm}(\mathbf{p}) \end{pmatrix} \\ V^{(\nu)}(\mathbf{p}, \pm\frac{1}{2}) &= \frac{1}{\sqrt{2}} \begin{pmatrix} \mp h_{\pm}(\mathbf{p}) \\ h_{\pm}(\mathbf{p}) \end{pmatrix} \end{aligned}$$

The four vectors $U^{(\nu)}(\mathbf{p}, \pm\frac{1}{2})$ and $V^{(\nu)}(\mathbf{p}, \pm\frac{1}{2})$ are an orthonormal basis in \mathbb{C}^4 .

Turning now to the theory of neutrinos and antineutrinos (see [16]) a neutrino has a helicity equal to $-\frac{1}{2}$ and an antineutrino a helicity equal to $\frac{1}{2}$. Neutrinos are left-handed and antineutrinos are right-handed. $U^{(\nu)}(\mathbf{p}, -\frac{1}{2})$ is the eigenfunction of a neutrino with a momentum \mathbf{p} and an energy $|\mathbf{p}|$. $CV^{(\nu)}(-\mathbf{p}, \frac{1}{2})$ is the eigenfunction of an antineutrino with a momentum \mathbf{p} and an energy $|\mathbf{p}|$.

Thus the classical field, denoted by $\Phi(\mathbf{x})$ and associated with the neutrino and the antineutrino, is given by

$$(3.25) \quad \begin{aligned} \Phi(\mathbf{x}) &= \left(\frac{1}{2\pi}\right)^{\frac{3}{2}} \text{L.i.m.} \left(\int_{\mathbb{R}^3} d^3\mathbf{p} \right. \\ &\quad \left. \left(e^{i(\mathbf{p}\cdot\mathbf{x})} U^{(\nu)}(\mathbf{p}, -\frac{1}{2}) a(\mathbf{p}, -\frac{1}{2}) + e^{-i(\mathbf{p}\cdot\mathbf{x})} V^{(\nu)}(-\mathbf{p}, \frac{1}{2}) c(\mathbf{p}, \frac{1}{2}) \right) \right). \end{aligned}$$

with

$$\|\Phi(\cdot)\|_{L^2(\mathbb{R}^3, \mathbb{C}^4)}^2 = \int_{\mathbb{R}^3} d^3p (|a(\mathbf{p}, -\frac{1}{2})|^2 + |c(\mathbf{p}, \frac{1}{2})|^2) < \infty.$$

3.2.1. Fock space for the neutrinos and the antineutrinos.

Let $\xi_4 = (\mathbf{p}, -\frac{1}{2})$ be the quantum variables of a neutrino where $\mathbf{p} \in \mathbb{R}^3$ is the momentum and $-\frac{1}{2}$ is the helicity. In the case of the antineutrino we set $\tilde{\xi}_4 = (\mathbf{p}, \frac{1}{2})$ where $\mathbf{p} \in \mathbb{R}^3$ and $\frac{1}{2}$ is the helicity.

$L^2(\mathbb{R}^3)$ is the Hilbert space of the states of the neutrinos and of the antineutrinos.

Let $\mathfrak{F}^{(\nu)}$ and $\mathfrak{F}^{(\bar{\nu})}$ denote the Fock spaces for the neutrinos and the antineutrinos respectively.

We have

$$(3.26) \quad \mathfrak{F}^{(\nu)} = \mathfrak{F}^{(\bar{\nu})} = \bigoplus_{n=0}^{\infty} \bigotimes_a^n L^2(\mathbb{R}^3)$$

$\bigotimes_a^n L^2(\mathbb{R}^3)$ is the antisymmetric n-th tensor power of $L^2(\mathbb{R}^3)$.
 $\Omega_{(\delta)} = (1, 0, 0, 0, \dots)$ is the vacuum state in $\mathfrak{F}^{(\delta)}$ for $\delta = \nu, \bar{\nu}$.

In the sequel we shall use the notations

$$(3.27) \quad \begin{aligned} \int_{\mathbb{R}^3} d\xi_4 &= \int_{\mathbb{R}^3} d^3p \\ \int_{\mathbb{R}^3} d\tilde{\xi}_4 &= \int_{\mathbb{R}^3} d^3p \end{aligned}$$

$b_+(\xi_4)$ (resp. $b_+^*(\xi_4)$) is the annihilation (resp. creation) operator for the neutrino and $b_-(\tilde{\xi}_4)$ (resp. $b_-^*(\tilde{\xi}_4)$) is the annihilation (resp. creation) operator for the antineutrino .

The operators $b_+(\xi_4)$, $b_+^*(\xi_4)$, $b_-(\tilde{\xi}_4)$ and $b_-^*(\tilde{\xi}_4)$ fulfil the usual anticommutation relations (CAR) and they anticommute with $b_\epsilon^\sharp(\xi_j)$ for $j = 1, 2, 3$ according the convention described in [28, Section 4.1]. See [6] arXiv for explicit definitions.

Therefore the following anticommutation relations hold for $j = 1, 2, 3$

$$(3.28) \quad \begin{aligned} \{b_+(\xi_4), b_+^*(\xi'_4)\} &= \delta(\xi_4 - \xi'_4) , \\ \{b_-(\tilde{\xi}_4), b_-^*(\tilde{\xi}'_4)\} &= \delta(\tilde{\xi}_4 - \tilde{\xi}'_4) , \\ \{b_+^\sharp(\xi_4), b_-^\sharp(\tilde{\xi}'_4)\} &= 0 , \\ \{b_+^\sharp(\xi_4), b_\epsilon^\sharp(\xi_j)\} &= \{b_-^\sharp(\tilde{\xi}_4), b_\epsilon^\sharp(\xi_j)\} = 0. \end{aligned}$$

Recall that for $\varphi \in L^2(\mathbb{R}^3)$, the operators

$$(3.29) \quad \begin{aligned} b_{4,+}(\varphi) &= \int_{\mathbb{R}^3} b_+(\xi_4) \overline{\varphi(\xi_4)} d\xi_4. \\ b_{4,-}(\varphi) &= \int_{\mathbb{R}^3} b_-(\tilde{\xi}_4) \overline{\varphi(\tilde{\xi}_4)} d\tilde{\xi}_4. \\ b_{4,+}^*(\varphi) &= \int_{\mathbb{R}^3} b_+^*(\xi_4) \varphi(\xi_4) d\xi_4. \\ b_{4,-}^*(\varphi) &= \int_{\mathbb{R}^3} b_-^*(\tilde{\xi}_4) \varphi(\tilde{\xi}_4) d\tilde{\xi}_4. \end{aligned}$$

are bounded operators on $\mathfrak{F}(\nu)$ and $\mathfrak{F}(\bar{\nu})$ respectively satisfying

$$(3.30) \quad \|b_{4,\epsilon}^\sharp(\varphi)\| = \|\varphi\|_{L^2}.$$

where $\epsilon = \pm$.

3.2.2. *Quantized Dirac Field for the neutrino.*

$e^{i(\mathbf{p}\cdot\mathbf{x})}U^{(\nu)}(\mathbf{p}, -\frac{1}{2})$ and $e^{i(\mathbf{p}\cdot\mathbf{x})}V^{(\nu)}(\mathbf{p}, \frac{1}{2})$ are generalized eigenfunctions of (3.21) with positive and negative eigenvalues $\pm E^{(\nu)}(\mathbf{p})$ respectively.

By (2.35) we get

$$(3.31) \quad C(V^{(\nu)}(\mathbf{p}, \frac{1}{2})) = \left(-\frac{p^1 + ip^2}{|p^1 + ip^2|}\right)U^{(\nu)}(-\mathbf{p}, \frac{1}{2})$$

Setting

$$(3.32) \quad \begin{aligned} U^{(\nu)}(\mathbf{p}, -\frac{1}{2}) &= U^{(\nu)}(\xi_4) \\ V^{(\nu)}(-\mathbf{p}, \frac{1}{2}) &= W^{(\nu)}(\tilde{\xi}_4). \end{aligned}$$

and applying the canonical quantization we obtain the following quantized Dirac field for the neutrino:

$$(3.33) \quad \begin{aligned} \Psi_{(\nu)}(\mathbf{x}) &= \left(\frac{1}{2\pi}\right)^{\frac{3}{2}} \left(\int d\xi_4 e^{i(\mathbf{p}\cdot\mathbf{x})} U^{(\nu)}(\xi_4) b_+(\xi_4) \right. \\ &\quad \left. + \int d\tilde{\xi}_4 e^{-i(\mathbf{p}\cdot\mathbf{x})} W^{(\nu)}(\tilde{\xi}_4) b_-^*(\tilde{\xi}_4) \right). \end{aligned}$$

4. THE HAMILTONIAN OF THE MODEL.

The processes (1.1) and (1.2) are associated with the β decay of the neutron (see [14],[16],[17] and [29]).

The β decay process can be described by the well known four-fermion effective Hamiltonian for the interaction in the Schrödinger representation:

$$(4.1) \quad \begin{aligned} H_{int} &= \\ &\frac{\tilde{G}}{\sqrt{2}} \int d^3x (\overline{\Psi}_{(p)}(x) \gamma^\alpha (1 - g_A \gamma_5) \Psi_{(n)}(x)) (\overline{\Psi}_{(e)}(x) \gamma_\alpha (1 - \gamma_5) \Psi_{(\nu)}(x)) \\ &+ \frac{\tilde{G}}{\sqrt{2}} \int d^3x (\overline{\Psi}_{(\nu)}(x) \gamma_\alpha (1 - \gamma_5) \Psi_{(e)}(x)) (\overline{\Psi}_{(n)}(x) \gamma^\alpha (1 - g_A \gamma_5) \Psi_{(p)}(x)) \end{aligned}$$

Here γ^α , $\alpha = 0, 1, 2, 3$ and γ_5 are the Dirac matrices in the standard representation. $\Psi_{(.)}(x)$ and $\overline{\Psi}_{(.)}(x)$ are the quantized Dirac fields for p , n , e and ν . $\overline{\Psi}_{(.)}(x) = \Psi_{(.)}^\dagger(x) \gamma^0$. $\tilde{G} = G_F \cos \theta_c$, where G_F is the Fermi coupling constant with $G_F \simeq 1.16639(2) \times 10^{-5} GeV^{-2}$ and θ_c is the Cabibbo angle with $\cos \theta_c \simeq 0.9751$. Moreover $g_A \simeq 1.27$. See [8].

The neutrino ν is the neutrino associated to the electron and usually denoted by ν_e in Physics.

From now on we restrict ourselves to the study of processes (1.1).

We recall that $m_e < m_p < m_p$.

4.1. The free Hamiltonian.

We set

$$\begin{aligned}
\mathfrak{F}^{(e)} &= \mathfrak{F}_{(e)} \otimes \mathfrak{F}_{(-e)}. \\
\mathfrak{F}^{(p)} &= \mathfrak{F}_{(p)} \otimes \mathfrak{F}_{(-p)}. \\
(4.2) \quad \mathfrak{F}^{(n)} &= \mathfrak{F}_{(n)} \\
\mathfrak{F}^{(\nu)} &= \mathfrak{F}_{(\nu)}. \\
\mathfrak{F} &= \mathfrak{F}^{(e)} \otimes \mathfrak{F}^{(p)} \otimes \mathfrak{F}^{(n)} \otimes \mathfrak{F}^{(\nu)}.
\end{aligned}$$

We set

$$\begin{aligned}
(4.3) \quad \omega(\xi_1) &= E_n^{(e)}(p^3) \quad \text{for } \xi_1 = (s, n, p^1, p^3) \\
\omega(\xi_2) &= E_n^{(p)}(p^3) \quad \text{for } \xi_2 = (s, n, p^1, p^3) \\
\omega(\xi_3) &= \sqrt{|\mathbf{p}|^2 + m_n^2} \quad \text{for } \xi_3 = (\mathbf{p}, \lambda) \\
\omega(\xi_4) &= |\mathbf{p}| \quad \text{for } \xi_4 = (\mathbf{p}, -\frac{1}{2}).
\end{aligned}$$

Let $H_D^{(e)}$ (resp. $H_D^{(p)}, H_D^{(n)}$ and $H_D^{(\nu)}$) be the Dirac Hamiltonian for the electron (resp. the proton, the neutron and the neutrino).

The quantization of $H_D^{(e)}$, denoted by $H_{0,D}^{(e)}$ and acting on $\mathfrak{F}^{(e)}$, is given by

$$(4.4) \quad H_{0,D}^{(e)} = \sum_{\epsilon=\pm} \int \omega(\xi_1) b_{\epsilon}^*(\xi_1) b_{\epsilon}(\xi_1) d\xi_1$$

Likewise the quantization of $H_D^{(p)}, H_D^{(n)}$ and $H_D^{(\nu)}$, denoted by $H_{0,D}^{(p)}, H_{0,D}^{(n)}$ and $H_{0,D}^{(\nu)}$ respectively, acting on $\mathfrak{F}^{(p)}, \mathfrak{F}^{(n)}$ and $\mathfrak{F}^{(\nu)}$ respectively, is given by

$$\begin{aligned}
(4.5) \quad H_{0,D}^{(p)} &= \sum_{\epsilon=\pm} \int \omega(\xi_2) b_{\epsilon}^*(\xi_2) b_{\epsilon}(\xi_2) d\xi_2 \\
H_{0,D}^{(n)} &= \int \omega(\xi_3) b_+^*(\xi_3) b_+(\xi_3) d\xi_3 \\
H_{0,D}^{(\nu)} &= \int \omega(\xi_4) b_+^*(\xi_4) b_+(\xi_4) d\xi_4.
\end{aligned}$$

For each Fock space $\mathfrak{F}^{(\cdot)}$, let $\mathfrak{D}^{(\cdot)}$ denote the set of vectors $\Phi \in \mathfrak{F}^{(\cdot)}$ for which each component $\Phi^{(r)}$ is smooth and has a compact support and $\Phi^{(r)} = 0$ for all but finitely many (r) . Then $H_{0,D}^{(\cdot)}$ is well-defined on the dense subset $\mathfrak{D}^{(\cdot)}$ and it is essentially self-adjoint on $\mathfrak{D}^{(\cdot)}$. The self-adjoint extension will be denoted by the same symbol $H_{0,D}^{(\cdot)}$ with domain $D(H_{0,D}^{(\cdot)})$.

The spectrum of $H_{0,D}^{(e)} \in \mathfrak{F}^{(e)}$ is given by

$$(4.6) \quad \text{spec}(H_{0,D}^{(e)}) = \{0\} \cup [m_e, \infty)$$

$\{0\}$ is a simple eigenvalue whose the associated eigenvector is the vacuum in $\mathfrak{F}^{(e)}$ denoted by $\Omega^{(e)}$. $[m_e, \infty)$ is the absolutely continuous spectrum of $H_{0,D}^{(e)}$.

Likewise the spectra of $H_{0,D}^{(p)}, H_{0,D}^{(n)}$ and $H_{0,D}^{(\nu)}$ are given by

$$(4.7) \quad \begin{aligned} \text{spec}(H_{0,D}^{(p)}) &= \{0\} \cup [m_p, \infty) \\ \text{spec}(H_{0,D}^{(n)}) &= \{0\} \cup [m_n, \infty) \\ \text{spec}(H_{0,D}^{(\nu)}) &= [0, \infty). \end{aligned}$$

$\Omega^{(p)}, \Omega^{(n)}$ and $\Omega^{(\nu)}$ are the associated vacua in $\mathfrak{F}^{(p)}, \mathfrak{F}^{(n)}$ and $\mathfrak{F}^{(\nu)}$ respectively and are the associated eigenvectors of $H_{0,D}^{(p)}, H_{0,D}^{(n)}$ and $H_{0,D}^{(\nu)}$ respectively for the eigenvalue $\{0\}$.

The vacuum in \mathfrak{F} , denoted by Ω , is then given by

$$(4.8) \quad \Omega = \Omega^{(e)} \otimes \Omega^{(p)} \otimes \Omega^{(n)} \otimes \Omega^{(\nu)}$$

The free Hamiltonian for the model, denoted by H_0 and acting on \mathfrak{F} , is now given by

$$(4.9) \quad \begin{aligned} H_0 &= H_{0,D}^{(e)} \otimes \mathbf{1} \otimes \mathbf{1} \otimes \mathbf{1} + \mathbf{1} \otimes H_{0,D}^{(p)} \otimes \mathbf{1} \otimes \mathbf{1} \\ &+ \mathbf{1} \otimes \mathbf{1} \otimes H_{0,D}^{(n)} \otimes \mathbf{1} + \mathbf{1} \otimes \mathbf{1} \otimes \mathbf{1} \otimes H_{0,D}^{(\nu)}. \end{aligned}$$

H_0 is essentially self-adjoint on $\mathfrak{D} = \mathfrak{D}^{(e)} \widehat{\otimes} \mathfrak{D}^{(p)} \widehat{\otimes} \mathfrak{D}^{(n)} \widehat{\otimes} \mathfrak{D}^{(\nu)}$.

Here $\widehat{\otimes}$ is the algebraic tensor product.

$\text{spec}(H_0) = [0, \infty)$ and Ω is the eigenvector associated with the simple eigenvalue $\{0\}$ of H_0 .

Let $S^{(e)}$ be the set of the thresholds of $H_{0,D}^{(e)}$:

$$S^{(e)} = (s_n^{(e)}; n \in \mathbb{N})$$

with $s_n^{(e)} = \sqrt{m_e^2 + 2neB}$.

Likewise let $S^{(p)}$ be the set of the thresholds of $H_{0,D}^{(p)}$:

$$S^{(p)} = (s_n^{(p)}; n \in \mathbb{N})$$

with $s_n^{(p)} = \sqrt{m_p^2 + 2neB}$.

Let $S^{(n)}$ be the set of the thresholds of $H_{0,D}^{(n)}$:

$$S^{(n)} = (nm_n; n \in \mathbb{N}, \text{ such that } n \geq 1)$$

Then

$$(4.10) \quad \mathfrak{S} = S^{(e)} \cup S^{(p)} \cup S^{(n)}$$

is the set of the thresholds of H_0 .

4.2. The Interaction.

By (4.1) let us now write down the formal interaction, denoted by V_I , involving the protons, the neutrons, the electrons and the neutrinos together with antiparticles in the Schrödinger representation for the process (1.1). We have

$$(4.11) \quad V_I = V_I^{(1)} + V_I^{(2)} + V_I^{(3)} + V_I^{(4)}$$

Set

$$(4.12) \quad \begin{aligned} \mathbf{q} &= \mathbf{p}_e + \mathbf{p}_p \\ \mathbf{r} &= \mathbf{p}_n + \mathbf{p}_\nu. \end{aligned}$$

After the integration with respect to (x^1, x^3) V_I is given by

$$(4.13) \quad \begin{aligned} V_I^{(1)} &= \int dx^2 \int d\xi_1 d\xi_2 d\xi_3 d\xi_4 e^{ix^2 r^2} \\ &(\overline{U^{(p)}}(x^2, \xi_2) \gamma^\alpha (1 - g_A \gamma_5) U^{(n)}(\xi_3)) (\overline{U^{(e)}}(x^2, \xi_1) \gamma_\alpha (1 - \gamma_5) U^{(\nu)}(\xi_4)) \\ &\delta(q^1 - r^1) \delta(q^3 - r^3) b_+^*(\xi_1) b_+^*(\xi_2) b_+(\xi_3) b_+(\xi_4). \end{aligned}$$

$$(4.14) \quad \begin{aligned} V_I^{(2)} &= \int dx^2 \int d\xi_1 d\xi_2 d\xi_3 d\xi_4 e^{-ix^2 r^2} \\ &(\overline{U^{(\nu)}}(\xi_4) \gamma_\alpha (1 - \gamma_5) U^{(e)}(x^2, \xi_1)) (\overline{U^{(n)}}(\xi_3) \gamma^\alpha (1 - g_A \gamma_5) U^{(p)}(x^2, \xi_2)) \\ &\delta(q^1 - r^1) \delta(q^3 - r^3) b_+^*(\xi_4) b_+^*(\xi_3) b_+(\xi_2) b_+(\xi_1). \end{aligned}$$

$$(4.15) \quad \begin{aligned} V_I^{(3)} &= \int dx^2 \int d\xi_1 d\xi_2 d\xi_3 d\xi_4 e^{-ix^2 r^2} \\ &(\overline{U^{(\nu)}}(\xi_4) \gamma_\alpha (1 - \gamma_5) W^{(e)}(x^2, \xi_1)) (\overline{U^{(n)}}(\xi_3) \gamma^\alpha (1 - g_A \gamma_5) W^{(p)}(x^2, \xi_2)) \\ &\delta(q^1 + r^1) \delta(q^3 + r^3) b_+^*(\xi_4) b_+^*(\xi_3) b_-(\xi_2) b_-(\xi_1). \end{aligned}$$

$$(4.16) \quad \begin{aligned} V_I^{(4)} &= \int dx^2 \int d\xi_1 d\xi_2 d\xi_3 d\xi_4 e^{ix^2 r^2} \\ &(\overline{W^{(p)}}(x^2, \xi_2) \gamma^\alpha (1 - g_A \gamma_5) U^{(n)}(\xi_3)) (\overline{W^{(e)}}(x^2, \xi_1) \gamma_\alpha (1 - \gamma_5) U^{(\nu)}(\xi_4)) \\ &\delta(q^1 + r^1) \delta(q^3 + r^3) b_+(\xi_4) b_+(\xi_3) b_-(\xi_2) b_-(\xi_1). \end{aligned}$$

$V_I^{(3)}$ and $V_I^{(4)}$ are responsible for the fact that the bare vacuum will not be an eigenvector of the total Hamiltonian as expected in Physics.

V_I is formally symmetric.

In the Fock space \mathfrak{F} the interaction V_I is a highly singular operator due to the δ -distributions that occur in the $(V_I^{(\cdot)})$'s and because of the ultraviolet behaviour of the functions $U^{(\cdot)}$ and $W^{(\cdot)}$.

In order to get well defined operators in \mathfrak{F} we have to substitute smoother kernels $F^{(\beta)}(\xi_2, \xi_3)$, $G^{(\beta)}(\xi_1, \xi_4)$, where $\beta = 1, 2$, both for the δ -distributions and the ultraviolet cutoffs .

We then obtain a new operator denoted by H_I and defined as follows in the Schrödinger representation.

$$(4.17) \quad H_I = H_I^{(1)} + H_I^{(2)} + H_I^{(3)} + H_I^{(4)}$$

with

$$(4.18) \quad \begin{aligned} H_I^{(1)} &= \int d\xi_1 d\xi_2 d\xi_3 d\xi_4 \left(\int dx^2 e^{ix^2 r^2} \right. \\ &\quad \left. \overline{(U^{(p)}(x^2, \xi_2) \gamma^\alpha (1 - g_A \gamma_5) U^{(n)}(\xi_3))} \overline{(U^{(e)}(x^2, \xi_1) \gamma_\alpha (1 - \gamma_5) U^{(\nu)}(\xi_4))} \right) \\ &\quad F^{(1)}(\xi_2, \xi_3) G^{(1)}(\xi_1, \xi_4) b_+^*(\xi_1) b_+^*(\xi_2) b_+(\xi_3) b_+(\xi_4). \end{aligned}$$

$$(4.19) \quad \begin{aligned} H_I^{(2)} &= \int d\xi_1 d\xi_2 d\xi_3 d\xi_4 \left(\int dx^2 e^{-ix^2 r^2} \right. \\ &\quad \left. \overline{(U^{(\nu)}(\xi_4) \gamma_\alpha (1 - \gamma_5) U^{(e)}(x^2, \xi_1))} \overline{(U^{(n)}(\xi_3) \gamma^\alpha (1 - g_A \gamma_5) U^{(p)}(x^2, \xi_2))} \right) \\ &\quad \overline{F^{(1)}(\xi_2, \xi_3) G^{(1)}(\xi_1, \xi_4) b_+^*(\xi_4) b_+^*(\xi_3) b_+(\xi_2) b_+(\xi_1)}. \end{aligned}$$

$$(4.20) \quad \begin{aligned} H_I^{(3)} &= \int d\xi_1 d\xi_2 d\xi_3 d\xi_4 \left(\int dx^2 e^{-ix^2 r^2} \right. \\ &\quad \left. \overline{(U^{(\nu)}(\xi_4) \gamma_\alpha (1 - \gamma_5) W^{(e)}(x^2, \xi_1))} \overline{(U^{(n)}(\xi_3) \gamma^\alpha (1 - g_A \gamma_5) W^{(p)}(x^2, \xi_2))} \right) \\ &\quad F^{(2)}(\xi_2, \xi_3) G^{(2)}(\xi_1, \xi_4) b_+^*(\xi_4) b_+^*(\xi_3) b_-^*(\xi_2) b_-^*(\xi_1). \end{aligned}$$

$$(4.21) \quad \begin{aligned} H_I^{(4)} &= \int d\xi_1 d\xi_2 d\xi_3 d\xi_4 \left(\int dx^2 e^{ix^2 r^2} \right. \\ &\quad \left. \overline{(W^{(p)}(x^2, \xi_2) \gamma^\alpha (1 - g_A \gamma_5) U^{(n)}(\xi_3))} \overline{(W^{(e)}(x^2, \xi_1) \gamma_\alpha (1 - \gamma_5) U^{(\nu)}(\xi_4))} \right) \\ &\quad \overline{F^{(2)}(\xi_2, \xi_3) G^{(2)}(\xi_1, \xi_4) b_+(\xi_4) b_+(\xi_3) b_-(\xi_2) b_-(\xi_1)}. \end{aligned}$$

Definition 4.1. *The total Hamiltonian is*

$$(4.22) \quad H = H_0 + gH_I$$

where g is a non-negative coupling constant.

The assumption that g is non-negative is made for simplicity but all the results below hold for $g \in \mathbb{R}$ with $|g|$ small enough.

We now give the hypothesis that the kernels $F^\beta(\cdot, \cdot), G^{(\beta)}(\cdot, \cdot)$, $\beta = 1, 2$, and the coupling constant g have to satisfy in order to associate with the formal operator H a well defined self-adjoint operator in \mathfrak{F} .

Throughout this work we assume the following hypothesis

Hypothesis 4.2. *For $\beta = 1, 2$ we assume*

$$(4.23) \quad \begin{aligned} F^{(\beta)}(\xi_2, \xi_3) &\in L^2(\Gamma_1 \times \Gamma_2) \\ G^{(\beta)}(\xi_1, \xi_4) &\in L^2(\Gamma_1 \times \mathbb{R}^3) \end{aligned}$$

Let $\langle \cdot, \cdot \rangle_{\mathbb{C}^4}$ be the scalar product in \mathbb{C}^4 . We have

$$(4.24) \quad \begin{aligned} \overline{U^{(p)}}(x^2, \xi_2) \gamma^\alpha (1 - g_A \gamma_5) U^{(n)}(\xi_3) &= \langle U^{(p)}(x^2, \xi_2), \gamma^0 \gamma^\alpha (1 - g_A \gamma_5) U^{(n)}(\xi_3) \rangle_{\mathbb{C}^4} \\ \overline{U^{(e)}}(x^2, \xi_1) \gamma_\alpha (1 - \gamma_5) U^{(\nu)}(\xi_4) &= \langle U^{(e)}(x^2, \xi_1), \gamma^0 \gamma_\alpha (1 - \gamma_5) U^{(\nu)}(\xi_4) \rangle_{\mathbb{C}^4} \\ \overline{U^{(\nu)}}(\xi_4) \gamma_\alpha (1 - \gamma_5) W^{(e)}(x^2, \xi_1) &= \langle U^{(\nu)}(\xi_4), \gamma^0 \gamma_\alpha (1 - \gamma_5) W^{(e)}(x^2, \xi_1) \rangle_{\mathbb{C}^4} \\ \overline{U^{(n)}}(\xi_3) \gamma^\alpha (1 - g_A \gamma_5) W^{(p)}(x^2, \xi_2) &= \langle U^{(n)}(\xi_3), \gamma^0 \gamma^\alpha (1 - g_A \gamma_5) W^{(p)}(x^2, \xi_2) \rangle_{\mathbb{C}^4}. \end{aligned}$$

Set

$$(4.25) \quad C_0 = \frac{1}{2} \left(\frac{1}{m_e} + \frac{1}{m_n} \right) (\|\gamma^\alpha (1 - g_A \gamma_5)\|) (\|\gamma_\alpha (1 - \gamma_5)\|)$$

We then have

Proposition 4.3. *For every $\Phi \in D(H_0)$ we obtain*

$$(4.26) \quad \begin{aligned} \|H_I^{(j)} \Phi\| &\leq C_0 \|F^{(1)}(\cdot, \cdot)\|_{L^2} \|G^{(1)}(\cdot, \cdot)\|_{L^2} \| (H_0 + m_n) \Phi \| \\ &\text{for } j=1, 2. \\ \|H_I^{(j)} \Phi\| &\leq C_0 \|F^{(2)}(\cdot, \cdot)\|_{L^2} \|G^{(2)}(\cdot, \cdot)\|_{L^2} \| (H_0 + m_n) \Phi \| \\ &\text{for } j=3, 4. \end{aligned}$$

By (4.23), (4.24) and (4.25) the estimates (4.26) are examples of N_τ estimates (see [15]). The proof is similar to the one of [5, Proposition 3.7] and details are omitted.

Let $g_0 > 0$ be such that

$$(4.27) \quad 2g_0 C_0 \left(\sum_{\beta=1}^2 \|F^{(\beta)}(\cdot, \cdot)\|_{L^2} \|G^{(\beta)}(\cdot, \cdot)\|_{L^2} \right) < 1$$

We now have

Theorem 4.4. *For any g such that $g \leq g_0$, H is a self-adjoint operator in \mathfrak{F} with domain $D(H) = D(H_0)$ and is bounded from below. H is essentially self-adjoint on any core of H_0 . Setting*

$$E = \inf \sigma(H)$$

we have for every $g \leq g_0$

$$\sigma(H) = \sigma_{\text{ess}}(H) = [E, \infty)$$

with $E \leq 0$.

Here $\sigma(H)$ is the spectrum of H and $\sigma_{\text{ess}}(H)$ is the essential spectrum of H .

Proof. By Proposition 4.2 and (4.27) the proof of the self-adjointness of H follows from the Kato-Rellich theorem.

We turn now to the essential spectrum. The result about the essential spectrum in the case of models involving bosons has been obtained by [12, theorem 4.1] and [2]. In the case of models involving fermions the result has been obtained by [26]. In our case involving only massive fermions and massless neutrinos we use the proof given by [26].

Thus we have to construct a Weyl sequence for H and $E + \lambda$ with $\lambda > 0$.

Let T be the self-adjoint multiplication operator in $L^2(\mathbb{R}^3)$ defined by $Tu(\mathbf{p}_4) = |\mathbf{p}_4|u(\mathbf{p}_4)$. T is the spectral representation of $H_D^{(\nu)}$ for the neutrinos of helicity $-\frac{1}{2}$ in the configuration space $L^2(\mathbb{R}^3)$. See (3.27).

Every $\lambda > 0$ belongs to the essential spectrum of T . Then there exists a Weyl sequence $(f_n)_{n \geq 1}$ for T and $\lambda > 0$ such that

$$(4.28) \quad \begin{aligned} f_n &\in D(T) \text{ for } n \geq 1. \\ \|f_n\| &= 1 \text{ for } n \geq 1. \\ w - \lim_{n \rightarrow \infty} f_n &= 0. \\ \lim_{n \rightarrow \infty} (T - \lambda)f_n &= 0. \end{aligned}$$

Let

$$(4.29) \quad \begin{aligned} f_n(\xi_4) &= f_n(\mathbf{p}_4) \\ b_{+,4}(f_n) &= \int b_+(\xi_4) \overline{f_n(\xi_4)} d\xi_4 \\ b_{+,4}^*(f_n) &= \int b_+^*(\xi_4) f_n(\xi_4) d\xi_4. \end{aligned}$$

In the following we identify $b_{+,4}^\sharp(f_n)$ with its obvious extension to \mathfrak{F} . An easy computation shows that, for every $\Psi \in D(H)$,

$$(4.30) \quad \begin{aligned} &\left[H_I^{(1)}, b_{+,4}^*(f_n) \right] \Psi = \\ &\int d\xi_1 d\xi_2 d\xi_3 \left(\int dx^2 e^{-ix^2 r^2} \right. \\ &\quad \left. \overline{U^{(p)}(x^2, \xi_2)} \gamma_\alpha (1 - g_A \gamma_5) U^{(ne)}(\xi_3) F^{(1)}(\xi_2, \xi_3) \right. \\ &\quad \left. \left\langle U^{(e)}(x^2, \xi_1), \gamma^0 \gamma^\alpha (1 - \gamma_5) \left(\int f_n(\xi_4) G^{(1)}(\xi_1, \xi_4) U^{(\nu_e)}(\xi_4) d\xi_4 \right) \right\rangle_{\mathbb{C}^4} \right) \\ &\quad b_+^*(\xi_1) b_+^*(\xi_2) b_+(\xi_3) \Psi. \end{aligned}$$

$$(4.31) \quad \left[H_I^{(1)}, b_{+,4}(f_n) \right] \Psi = 0$$

$$(4.32) \quad \begin{aligned} & \left[H_I^{(2)}, b_{+,4}(f_n) \right] \Psi = \\ & - \int d\xi_1 d\xi_2 d\xi_3 \left(\int dx^2 e^{-ix^2 r^2} \right. \\ & \left. \overline{(U^{(ne)}(\xi_3) \gamma^\alpha (1 - g_A \gamma_5) U^{(p)}(x^2, \xi_2)) F^{(1)}(\xi_2, \xi_3)} \right. \\ & \left. \left\langle \int f_n(\xi_4) G^{(1)}(\xi_1, \xi_4) U^{(\nu)}(\xi_4) d\xi_4, \gamma^0 \gamma^\alpha (1 - \gamma_5) U^{(e)}(x^2, \xi_1) \right\rangle_{\mathbb{C}^4} \right) \\ & b_+^*(\xi_3) b_+(\xi_2) b_+(\xi_1) \Psi. \end{aligned}$$

$$(4.33) \quad \left[H_I^{(2)}, b_{+,4}^*(f_n) \right] \Psi = 0$$

$$(4.34) \quad \begin{aligned} & \left[H_I^{(3)}, b_{+,4}(f_n) \right] \Psi = \\ & - \int d\xi_1 d\xi_2 d\xi_3 \left(\int dx^2 e^{-ix^2 r^2} \right. \\ & \left. \overline{(U^{(ne)}(\xi_3) \gamma^\alpha (1 - g_A \gamma_5) W^{(p)}(x^2, \xi_2)) F^{(2)}(\xi_2, \xi_3)} \right. \\ & \left. \left\langle \int f_n(\xi_4) \overline{G^{(2)}(\xi_1, \xi_4) U^{(\nu)}(\xi_4) d\xi_4}, \gamma_0 \gamma^\alpha (1 - \gamma_5) W^{(e)}(x^2, \xi_1) \right\rangle_{\mathbb{C}^4} \right) \\ & b_+^*(\xi_3) b_-^*(\xi_2) b_-^*(\xi_1) \Psi. \end{aligned}$$

$$(4.35) \quad \left[H_I^{(3)}, b_{+,4}^*(f_n) \right] \Psi = 0$$

$$(4.36) \quad \begin{aligned} & \left[H_I^{(4)}, b_{+,4}^*(f_n) \right] \Psi = \\ & \int d\xi_1 d\xi_2 d\xi_3 \left(\int dx^2 e^{-ix^2 r^2} \right. \\ & \left. \overline{(W^{(p)}(x^2, \xi_2) \gamma^\alpha (1 - g_A \gamma_5) U^{(ne)}(\xi_3)) F^{(2)}(\xi_2, \xi_3)} \right. \\ & \left. \left\langle W^{(e)}(x^2, \xi_1), \gamma^0 \gamma^\alpha (1 - \gamma_5) \left(\int f_n(\xi_4) \overline{G^{(2)}(\xi_1, \xi_4) U^{(\nu)}(\xi_4) d\xi_4} \right) \right\rangle_{\mathbb{C}^4} \right) \\ & b_+(\xi_3) b_-(\xi_2) b_-(\xi_1) \Psi. \end{aligned}$$

$$(4.37) \quad \left[H_I^{(4)}, b_{+,4}(f_n) \right] \Psi = 0$$

Let $P_H(\cdot)$ be the spectral measure of H . For any $\epsilon > 0$ the orthogonal projection $P_H([E, E + \epsilon])$ is different from zero because E belongs to $\sigma(H)$.

Let $\Phi_\epsilon \in \text{Ran}(P_H([E, E + \epsilon]))$ such that $\|\Phi_\epsilon\| = 1$. We set

$$(4.38) \quad \Psi_{n,\epsilon} = (b_{+,4}(f_n) + b_{+,4}^*(f_n))\Phi_\epsilon, \quad n \geq 1$$

Let us show that there exists a subsequence of $(\Psi_{n,\epsilon})_{n \geq 1, \epsilon > 0}$ which is a Weyl sequence for H and $E + \lambda$ with $\lambda > 0$.

By Hypothesis 4.1, (4.30), (4.32), (4.34), (4.36) and the N_τ estimates we get

$$(4.39) \quad \begin{aligned} & \sup \left(\left\| \left[H_I^{(1)}, b_{+,4}(f_n) \right] \Psi \right\|, \left\| \left[H_I^{(2)}, b_{+,4}(f_n) \right] \Psi \right\| \right) \\ & \leq C_0 \|F^{(1)}(\cdot, \cdot)\|_{L^2(\Gamma_1 \times \Gamma_2)} \left(\int \left\| \int f_n(\xi_4) G^{(1)}(\xi_1, \xi_4) U^{(\nu)}(\xi_4) d\xi_4 \right\|_{\mathbb{C}^4}^2 d\xi_1 \right)^{\frac{1}{2}} \\ & \qquad \qquad \qquad \| (H_0 + m_p)^{\frac{1}{2}} \Psi \|. \\ & \sup \left(\left\| \left[H_I^{(3)}, b_{+,4}(f_n) \right] \Psi \right\|, \left\| \left[H_I^{(4)}, b_{+,4}(f_n) \right] \Psi \right\| \right) \\ & \leq C_0 \|F^{(2)}(\cdot, \cdot)\|_{L^2(\Gamma_1 \times \Gamma_2)} \left(\int \left\| \int f_n(\xi_4) \overline{G^{(2)}(\xi_1, \xi_4)} U^{(\nu)}(\xi_4) d\xi_4 \right\|_{\mathbb{C}^4}^2 d\xi_1 \right)^{\frac{1}{2}} \\ & \qquad \qquad \qquad \| (H_0 + m_p)^{\frac{1}{2}} \Psi \|. \end{aligned}$$

Note that

$$(4.40) \quad \|\Psi_{n,\epsilon}\| = 1, \quad n \geq 1$$

We have for every $\Psi \in D(H)$

$$(4.41) \quad \begin{aligned} & (H\Psi, \Psi_{n,\epsilon}) = \\ & (\Psi, (b_{+,4}(f_n) + b_{+,4}^*(f_n))H\Phi_\epsilon + (b_{+,4}^*(Tf_n) - (b_{+,4}(Tf_n)))\Phi_\epsilon \\ & \quad + g[H_I, (b_{+,4}(f_n) + b_{+,4}^*(f_n))]\Psi_\epsilon). \end{aligned}$$

See [12].

This yields

$$(4.42) \quad \begin{aligned} & H\Psi_{n,\epsilon} = \\ & ((b_{+,4}(f_n) + b_{+,4}^*(f_n))H\Phi_\epsilon + (b_{+,4}^*(Tf_n) - (b_{+,4}(Tf_n)))\Phi_\epsilon \\ & \quad + g[H_I, (b_{+,4}(f_n) + b_{+,4}^*(f_n))]\Psi_\epsilon). \end{aligned}$$

and

$$\begin{aligned}
(4.43) \quad & (H - E - \lambda)\Psi_{n,\epsilon} = \\
& (b_{+,4}(f_n) + b_{+,4}^*(f_n))(H - E)\Psi_\epsilon \\
& + (b_{+,4}((T + \lambda)f_n) + b_{+,4}^*((T - \lambda)f_n))\Psi_\epsilon \\
& + g[H_I, (b_{+,4}(f_n) + b_{+,4}^*(f_n))]\Psi_\epsilon.
\end{aligned}$$

By (3.19) this yields for $|g| \leq g_0$

$$\begin{aligned}
(4.44) \quad & \|(H - E - \lambda)\Psi_{n,\epsilon}\| \leq \\
& + 2\epsilon + 2|\lambda|\|b_{+,4}(f_n)\Psi_\epsilon\| + 2\|((T - \lambda)f_n)\| \\
& + |g|\|[H_I, b_{+,4}(f_n)]\Psi_\epsilon\| + |g|\|[H_I, b_{+,4}^*(f_n)]\Psi_\epsilon\|
\end{aligned}$$

Let $\{g_k | k = 1, 2, 3, \dots\}$ be an orthonormal basis of $L^2(\mathbb{R}^3)$ and consider

$$(4.45) \quad b_{+,4}^*(g_{k_1})b_{+,4}^*(g_{k_2})b_{+,4}^*(g_{k_3})\dots\dots b_{+,4}^*(g_{k_m})\Omega_\nu \in \mathfrak{F}(\nu)$$

where the indices can be assumed ordered $k_1 < \dots < k_m$. Fock space vectors of this type form a basis of $\mathfrak{F}(\nu)$ (see [27]). By [26, Lemma 2.1] this yields for every $\epsilon > 0$

$$\begin{aligned}
(4.46) \quad & s - \lim_{n \rightarrow \infty} b_{+,4}(f_n)\Psi_\epsilon = 0, \\
& w - \lim_{n \rightarrow \infty} b_{+,4}^*(f_n)\Psi_\epsilon = 0.
\end{aligned}$$

By (3.26) and Hypothesis 4.1 we have

$$\begin{aligned}
(4.47) \quad & \lim_{n \rightarrow \infty} \left(\int \left\| \int f_n(\xi_4)G^{(1)}(\xi_1, \xi_4)U^{(\nu)}(\xi_4)d\xi_4 \right\|_{\mathbb{C}^4}^2 d\xi_1 \right)^{\frac{1}{2}} = 0 \\
& \lim_{n \rightarrow \infty} \left(\int \left\| \int f_n(\xi_4)\overline{G^{(2)}(\xi_1, \xi_4)}U^{(\nu)}(\xi_4)d\xi_4 \right\|_{\mathbb{C}^4}^2 d\xi_1 \right)^{\frac{1}{2}} = 0.
\end{aligned}$$

It follows from (4.28), (4.38), (4.44), (4.46) and (4.47) that for every $\epsilon > 0$

$$(4.48) \quad \limsup_{n \rightarrow \infty} \|(H - E - \lambda)\Psi_{n,\epsilon}\| \leq 2\epsilon$$

This yields

$$(4.49) \quad \lim_{\epsilon \rightarrow 0} \limsup_{n \rightarrow \infty} \|(H - E - \lambda)\Psi_{n,\epsilon}\| = 0$$

In view of (4.49) there exists a subsequence $(\Psi_{n_j, \epsilon_j})_{j \geq 1}$ such that

$$(4.50) \quad \lim_{j \rightarrow \infty} \|(H - E - \lambda)\Psi_{n_j, \epsilon_j}\| = 0$$

Furthermore it follows from (4.46) that $w - \lim_{j \rightarrow \infty} \Psi_{n_j, \epsilon_j} = 0$.

The sequence $(\Psi_{n_j, \epsilon_j})_{j \geq 1}$ is a Weyl sequence for H and $E + \lambda$ with $\lambda > 0$

In order to show that $E \leq 0$ we adapt the proof given in [4] and [7]. We omit the details.

This concludes the proof of theorem 4.4. □

5. EXISTENCE OF A UNIQUE GROUND STATE FOR THE HAMILTONIAN H .

Set

$$(5.1) \quad \begin{aligned} K(F, G) &= \sum_{\beta=1}^2 \|F^{(\beta)}(\cdot, \cdot)\|_{L^2} \|G^{(\beta)}(\cdot, \cdot)\|_{L^2}. \\ C &= 2C_0. \\ B &= 2m_n C_0. \end{aligned}$$

By (4.26) and (5.1) we get for every $\psi \in D(H)$

$$(5.2) \quad \|H_T \psi\| \leq K(F, G) (C \|H_0 \psi\| + B \|\psi\|)$$

In order to prove the existence of a ground state for the Hamiltonian H we shall make the following additional assumptions on the kernels $G^{(\beta)}(\xi_1, \xi_4)$, $\beta = 1, 2$.

From now on $\mathbf{p}_4 \in \mathbb{R}^3$ is the momentum of the neutrino with helicity $-\frac{1}{2}$.

Hypothesis 5.1. *There exists a constant $\tilde{K}(G) > 0$ such that for $\beta = 1, 2$ and $\sigma > 0$*

$$(i) \quad \int_{\Gamma_1 \times \mathbb{R}^3} \frac{|G^{(\beta)}(\xi_1, \xi_4)|^2}{|\mathbf{p}_4|^2} d\xi_1 d\xi_4 < \infty ,$$

$$(ii) \quad \left(\int_{\Gamma_1 \times \{|\mathbf{p}_4| \leq \sigma\}} |G^{(\beta)}(\xi_1, \xi_4)|^2 d\xi_1 d\xi_4 \right)^{\frac{1}{2}} \leq \tilde{K}(G) \sigma.$$

We have

Theorem 5.2. *Assume that the kernels $F^{(\beta)}(\cdot, \cdot)$ and $G^{(\beta)}(\cdot, \cdot)$, $\beta = 1, 2$, satisfy Hypothesis 4.1 and Hypothesis 5.1. Then there exists $g_1 \in (0, g_0]$ such that H has a unique ground state for $g \leq g_1$.*

In order to prove theorem 5.2 we first prove the existence of a spectral gap for some neutrino infrared cutoff Hamiltonians.

5.1. The neutrino infrared cutoff Hamiltonians and the existence of a spectral gap.

Proof. Let us first define the neutrino infrared cutoff Hamiltonians.

For that purpose, let $\chi_0(\cdot) \in C^\infty(\mathbb{R}, [0, 1])$ with $\chi_0 = 1$ on $(-\infty, 1]$ and $\chi_0 = 0$ on $[2, \infty)$. For $\sigma > 0$ and $\mathbf{p}_4 \in \mathbb{R}^3$, we set

$$(5.3) \quad \begin{aligned} \chi_\sigma(\mathbf{p}_4) &= \chi_0(|\mathbf{p}_4|/\sigma) , \\ \tilde{\chi}^\sigma(\mathbf{p}_4) &= 1 - \chi_\sigma(\mathbf{p}_4) . \end{aligned}$$

The operator $H_{I,\sigma}$ is the interaction given by (4.17) associated with the kernels $F^{(\beta)}(\xi_2, \xi_3)\tilde{\chi}^\sigma(\mathbf{p}_4)G^{(\beta)}(\xi_1, \xi_4)$ instead of $F^{(\beta)}(\xi_2, \xi_3)G^{(\beta)}(\xi_1, \xi_4)$.

We then set

$$(5.4) \quad H_\sigma = H_0 + gH_{I,\sigma} .$$

We now introduce

$$(5.5) \quad \begin{aligned} \Gamma_{4,\sigma} &= \mathbb{R}^3 \cap \{|\mathbf{p}_4| < \sigma\} , & \Gamma_4^\sigma &= \mathbb{R}^3 \cap \{|\mathbf{p}_4| \geq \sigma\} \\ \mathfrak{F}_{4,\sigma} &= \mathfrak{F}_a(L^2(\Gamma_{4,\sigma})) , & \mathfrak{F}_4^\sigma &= \mathfrak{F}_a(L^2(\Gamma_4^\sigma)) . \end{aligned}$$

$\mathfrak{F}_{4,\sigma} \otimes \mathfrak{F}_4^\sigma$ is the Fock space for the massless neutrino such that $\mathfrak{F}^{(\nu)} \simeq \mathfrak{F}_{4,\sigma} \otimes \mathfrak{F}_4^\sigma$.

We set

$$(5.6) \quad \mathfrak{F}^\sigma = \mathfrak{F}^{(e)} \otimes \mathfrak{F}^{(p)} \otimes \mathfrak{F}^{(n)} \otimes \mathfrak{F}_4^\sigma \text{ and } \mathfrak{F}_\sigma = \mathfrak{F}_{4,\sigma}$$

We have

$$(5.7) \quad \mathfrak{F} \simeq \mathfrak{F}^\sigma \otimes \mathfrak{F}_\sigma .$$

We further set

$$(5.8) \quad H_0^4 = \int |\mathbf{p}_4| b_+^*(\xi_4) b_+(\xi_4) d\xi_4 .$$

In the following we identify H_0^4 with its obvious extension to \mathfrak{F} .

We let

$$(5.9) \quad \begin{aligned} H_0^{4,\sigma} &= \int_{|\mathbf{p}_4| \geq \sigma} |\mathbf{p}_4| b_+^*(\xi_4) b_+(\xi_4) d\xi_4 , \\ H_{0,\sigma}^4 &= \int_{|\mathbf{p}_4| < \sigma} |\mathbf{p}_4| b_+^*(\xi_4) b_+(\xi_4) d\xi_4 . \end{aligned}$$

We identify $H_0^{4,\sigma}$ and $H_{0,\sigma}^4$ with their obvious extension to \mathfrak{F}^σ and \mathfrak{F}_σ respectively.

On $\mathfrak{F}^\sigma \otimes \mathfrak{F}_\sigma$, we have

$$(5.10) \quad H_0^4 = H_0^{4,\sigma} \otimes \mathbf{1}_\sigma + \mathbf{1}^\sigma \otimes H_{0,\sigma}^4 .$$

where $\mathbf{1}^\sigma$ (resp. $\mathbf{1}_\sigma$) is the identity operator on \mathfrak{F}^σ (resp. \mathfrak{F}_σ).

Setting

$$(5.11) \quad H_0^\sigma = H_0|_{\mathfrak{F}^\sigma} \text{ and } H^\sigma = H_\sigma|_{\mathfrak{F}^\sigma} .$$

we then get

$$(5.12) \quad \begin{aligned} H_0^\sigma &= H_{0,D}^{(e)} + H_{0,D}^{(p)} + H_{0,D}^{(n)} + H_0^{4,\sigma} \text{ on } \mathfrak{F}^\sigma , \\ H^\sigma &= H_0^\sigma + gH_{I,\sigma} \text{ on } \mathfrak{F}^\sigma . \end{aligned}$$

and

$$(5.13) \quad H_\sigma = H^\sigma \otimes \mathbf{1}_\sigma + \mathbf{1}^\sigma \otimes H_{0,\sigma}^4 . \text{ on } \mathfrak{F}^\sigma \otimes \mathfrak{F}_\sigma .$$

On the other hand, for $\delta \in \mathbb{R}$ such that $0 < \delta < m_3$, we define the sequence $(\sigma_n)_{n \geq 0}$ by

$$(5.14) \quad \begin{aligned} \sigma_0 &= 2m_e + 1, \\ \sigma_1 &= m_e - \frac{\delta}{2}, \\ \sigma_{n+1} &= \gamma \sigma_n \text{ for } n \geq 1, \end{aligned}$$

where

$$(5.15) \quad \gamma = 1 - \frac{\delta}{2m_e - \delta}.$$

For $n \geq 0$, we now introduce the neutrino infrared cutoff Hamiltonians on $\mathfrak{F}^n = \mathfrak{F}^{\sigma_n}$ by setting

$$(5.16) \quad H^n = H^{\sigma_n}, \quad H_0^n = H_0^{\sigma_n},$$

We set, for $n \geq 0$,

$$(5.17) \quad E^n = \inf \sigma(H^n).$$

We introduce the neutrino infrared cutoff Hamiltonians on \mathfrak{F} by setting

$$(5.18) \quad H_n = H_{\sigma_n}, \quad H_{0,n} = H_{0,\sigma_n},$$

We set, for $n \geq 0$,

$$(5.19) \quad E_n = \inf \sigma(H_n).$$

Note that

$$(5.20) \quad E^n = E_n$$

One easily shows that, for $g \leq g_0$,

$$(5.21) \quad |E^n| = |E_n| \leq \frac{gK(F, G)B}{1 - g_0K(F, G)C}$$

See [6, 4] for a proof.

We now let

$$(5.22) \quad \tilde{K}(F, G) = 2 \left(\sum_{\beta=1,2} \|F^{(\beta)}(\cdot, \cdot)\|_{L^2(\Gamma_1 \times \Gamma_1)} \right) \tilde{K}(G)$$

where $\tilde{K}(G)$ is the constant given in Hypothesis 5.2(ii).

We further set,

$$(5.23) \quad \tilde{C} = \frac{C}{(1 - g_0K(F, G)C)},$$

$$(5.24) \quad \tilde{B} = \frac{B}{(1 - g_0K(F, G)C)^2}$$

and

$$(5.25) \quad \tilde{D}(F, G) = \max \left\{ \frac{4(2m_3 + 1)\gamma}{2m_3 - \delta}, 2 \right\} \tilde{K}(F, G)(2m_3\tilde{C} + \tilde{B}).$$

Let $g_1^{(\delta)}$ be such that

$$(5.26) \quad 0 < g_1^{(\delta)} < \min \left\{ 1, g_0, \frac{\gamma - \gamma^2}{3\tilde{D}(F, G)} \right\}.$$

and let

$$(5.27) \quad g_3 = \frac{1}{2K(F, G)(2C + B)}$$

Setting

$$(5.28) \quad \begin{aligned} g_2^{(\delta)} &= \inf \{ g_3, g_1^{(\delta)} \} \\ C(F, G) &= 3 \frac{\tilde{D}(F, G)}{\gamma}. \end{aligned}$$

and applying the same method as the one used for proving proposition 4.1 in [4] we finally get the existence of a spectral gap for H^n . We omit the details of the proof.

The proof of the following proposition is achieved.

Proposition 5.3. *Suppose that the kernels $F^{(\beta)}(., .), G^{(\beta)}(., .), \beta = 1, 2$, satisfy Hypothesis 4.1 and Hypothesis 5.1(ii). Then, for $g \leq g_2^{(\delta)}$, E^n is a simple eigenvalue of H^n for $n \geq 1$, and H^n does not have spectrum in the interval $(E^n, E^n + (1 - gC(F, G))\sigma_n)$.*

□

5.2. Proof of the existence of a ground state.

Proof. In order to prove the existence of a ground state for H we adapt the proof of theorem 3.3 in [6]. By Proposition 5.3 H^n has a unique ground state, denoted by ϕ^n , in \mathfrak{F}^n such that

$$(5.29) \quad H^n \phi^n = E^n \phi^n, \quad \phi^n \in \mathcal{D}(H^n), \quad \|\phi^n\| = 1, \quad n \geq 1.$$

Therefore H_n has a unique normalized ground state in \mathfrak{F} , given by $\tilde{\phi}_n = \phi^n \otimes \Omega_n$, where Ω_n is the vacuum state in \mathfrak{F}_n ,

$$(5.30) \quad H_n \tilde{\phi}_n = E^n \tilde{\phi}_n, \quad \tilde{\phi}_n \in \mathcal{D}(H_n), \quad \|\tilde{\phi}_n\| = 1, \quad n \geq 1.$$

Let $H_{I,n}$ denote the interaction H_{I,σ_n} . It follows from the pull-through formula that

$$(5.31) \quad (H_0 + gH_{I,n})b_+(\xi_4)\tilde{\phi}_n = E_n b_+(\xi_4)\tilde{\phi}_n - \omega(\xi_4)b_+(\xi_4)\tilde{\phi}_n - (g\tilde{V}_n^1(\xi_4) + g\tilde{V}_n^2(\xi_4))\tilde{\phi}_n$$

where

(5.32)

$$\begin{aligned} \tilde{V}_n^{(1)}(\xi_4) &= \int d\xi_1 d\xi_2 d\xi_3 \left(\int dx^2 e^{-ix^2 r^2} \right. \\ &\quad \left. \overline{U^{(\nu_e)}(\xi_4) \gamma_\alpha (1 - \gamma_5) U^{(e)}(x^2, \xi_1)} \overline{U^{(ne)}(\xi_3) \gamma^\alpha (1 - g_A \gamma_5) U^{(p)}(x^2, \xi_2)} \right) \\ &\quad \overline{F^{(1)}(\xi_2, \xi_3) G^{(1)}(\xi_1, \xi_4) \tilde{\chi}^{\sigma_n}(\mathbf{p}_4) b_+^*(\xi_3) b_+(\xi_2) b_+(\xi_1)}. \end{aligned}$$

(5.33)

$$\begin{aligned} \tilde{V}_n^{(2)}(\xi_4) &= \int d\xi_1 d\xi_2 d\xi_3 \left(\int dx^2 e^{-ix^2 r^2} \right. \\ &\quad \left. \overline{U^{(\nu_e)}(\xi_4) \gamma_\alpha (1 - \gamma_5) W^{(e)}(x^2, \xi_1)} \overline{U^{(ne)}(\xi_3) \gamma^\alpha (1 - g_A \gamma_5) W^{(p)}(x^2, \xi_2)} \right) \\ &\quad F^{(2)}(\xi_2, \xi_3) G^{(2)}(\xi_1, \xi_4) \tilde{\chi}^{\sigma_n}(\mathbf{p}_4) b_+^*(\xi_3) b_-^*(\xi_2) b_-^*(\xi_1). \end{aligned}$$

Hence, by (5.30), (5.31), (5.32) and (5.33), we get

$$(5.34) \quad (H_n - E_n + \omega(\xi_4)) b_+(\xi_4) \tilde{\phi}_n = -g \left(\tilde{V}_n^{(1)}(\xi_4) + \tilde{V}_n^{(2)}(\xi_4) \right) \tilde{\phi}_n.$$

We further note that, for $\beta = 1, 2,$,

$$(5.35) \quad \|\tilde{V}_n^{(\beta)}(\xi_4) \tilde{\phi}_n\| \leq \tilde{C}_0 \|F^{(\beta)}(\cdot, \cdot)\|_{L^2(\Gamma_1 \times \Gamma_1)} \|G^{(\beta)}(\cdot, \xi_4)\|_{L^2(\Gamma_1)} \times \|(H_0 + m_n)^{\frac{1}{2}} \tilde{\phi}_n\|.$$

where

$$\tilde{C}_0 = \left(\frac{1}{m_n}\right)^{\frac{1}{2}} (\|\gamma^\alpha (1 - g_A \gamma_5)\|) (\|\gamma_\alpha (1 - \gamma_5)\|)$$

The estimates (5.35) are examples of N_τ estimates (see [15]). The proof is similar to the one of [5, Proposition 3.7] and details are omitted.

Let us estimate $\|H_0 \tilde{\phi}_n\|$. By (5.2) we get

$$(5.36) \quad g \|H_{I,n} \tilde{\phi}_n\| \leq gK(F, G) \left(C \|H_0 \tilde{\phi}_n\| + B \right)$$

and

$$(5.37) \quad \|H_0 \tilde{\phi}_n\| \leq |E_n| + |g| \|gH_{I,n} \tilde{\phi}_n\|$$

By (5.21), we obtain

$$(5.38) \quad \|H_0 \tilde{\phi}_n\| \leq \frac{g_0 K(F, G) B}{1 - g_0 K(F, G) C} \left(1 + \frac{1}{1 - g_0 K(F, G) C} \right) = M$$

By (5.38) $\|H_0 \tilde{\phi}_n\|$ is bounded uniformly with respect to n and $g \leq g_0$ and by (5.34), (5.35) and (5.38) we get

$$(5.39) \quad \|b_+(\xi_4)\tilde{\phi}_n\| \leq \frac{g\tilde{C}_0}{|\mathbf{P}_4|} \left(\sum_{\beta=1}^2 \|F^{(\beta)}(\cdot, \cdot)\|_{L^2} \|G^{(\beta)}(\cdot, \xi_4)\|_{L^2} \right) (M + m_p)^{\frac{1}{2}}$$

uniformly with respect to n .

By Hypothesis 5.1(i) and (5.39) there exists a constant $\tilde{C}(F, G) > 0$ such that

$$(5.40) \quad \int \|b_+(\xi_4)\tilde{\phi}_n\|^2 d\xi_4 \leq \tilde{C}(F, G)g^2$$

Since $\|\tilde{\phi}_n\| = 1$, there exists a subsequence $(n_k)_{k \geq 1}$, converging to ∞ such that $(\tilde{\phi}_{n_k})_{k \geq 1}$ converges weakly to a state $\tilde{\phi} \in \mathfrak{F}$. By adapting the proof of theorem 4.1 in [5, 1] it follows from (5.40) that there exists g_2 such that $0 < g_2 \leq g_2^{(\delta)}$ and $\tilde{\phi} \neq 0$ for any $g \leq g_2$. Thus $\tilde{\phi}$ is a ground state of H . \square

5.3. Uniqueness of a ground state of the Hamiltonian H .

Proof. The proof follows by adapting the one given in [7]. See also [3].

In view of theorem 4.3 E is an eigenvalue of H with a finite multiplicity. Either E is a simple eigenvalue and the theorem is proved or its multiplicity is equal to $p \in \mathbb{N}$ with $p > 1$. Let us consider the second case. We wish to show by contradiction that E is a simple eigenvalue for g sufficiently small.

Let (ϕ_1, ϕ_2) be two vectors of the eigenspace of E . Each ϕ_j with $j = 1, 2$ is a ground state of H . ϕ_1 and ϕ_2 can be chosen such that $\langle \phi_1, \phi_2 \rangle_{\mathfrak{F}} = 0$ with $\|\phi_j\| = 1$, $j = 1, 2$.

By (5.30) let ϕ_n be a unique normalized ground state of H_n .

We have

$$(5.41) \quad \begin{aligned} 0 &= |\langle \phi_2, \phi_1 \rangle|^2 = \lim_{n \rightarrow \infty} |\langle \phi_2, \phi_n \rangle|^2 \\ &= \lim_{n \rightarrow \infty} \langle E_{\{E_n\}}(H_n)\phi_2, \phi_2 \rangle \\ &= 1 - \lim_{n \rightarrow \infty} \langle (\mathbf{1}_{\mathfrak{F}} - E_{\{E_n\}}(H_n))\phi_2, \phi_2 \rangle. \end{aligned}$$

where $E_{\{\cdot\}}(\cdot)$ is the spectral measure for the associated self-adjoint operator.

We have

$$(5.42) \quad \mathbf{1}_{\mathfrak{F}} - E_{\{E_n\}}(H_n) = (\mathbf{1}^n - E_{\{E_n\}}(H^n)) \otimes P_{\Omega_n} + \mathbf{1}^n \otimes (\mathbf{1}_n - P_{\Omega_n})$$

We have to estimate

$$(5.43) \quad \langle \phi_2, (\mathbf{1}^n \otimes (\mathbf{1}_n - P_{\Omega_n}))\phi_2 \rangle,$$

and

$$(5.44) \quad \langle \phi_2, ((\mathbf{1}^n - E_{\{E_n\}}(H^n)) \otimes P_{\Omega_n})\phi_2 \rangle$$

We first estimate (5.43).

By applying the same proof as the one used to get estimates (5.38), (5.39) and (5.40) with ϕ_2 instead of $\tilde{\phi}_n$ we easily get

$$(5.45) \quad \int \|b_+(\xi_4)\phi_2\|^2 d\xi_4 \leq \tilde{C}(F, G)g^2$$

This yields

$$(5.46) \quad \langle \phi_2, (\mathbf{1}^n \otimes (\mathbf{1}_n - P_{\Omega_n})\phi_2) \rangle \leq \tilde{C}(F, G)g^2$$

We now estimate (5.44)

Set

$$(5.47) \quad \mathbf{1}^n - E_{\{E_n\}}(H^n) = E_{\{E_n\}}(H^n)^\perp$$

By proposition 5.3 we get

$$(5.48) \quad (H^n - E_n)E_{\{E_n\}}(H^n)^\perp \geq (1 - gC(F, G))\sigma_n E_{\{E_n\}}(H^n)^\perp.$$

and

$$(5.49) \quad \begin{aligned} & \langle \phi_2, E_{\{E_n\}}(H^n)^\perp \otimes P_{\Omega_n}\phi_2 \rangle \\ & \leq \frac{1}{(1 - gC(F, G))\sigma_n} \langle \phi_2, (H^n - E_n) \otimes P_{\Omega_n}\phi_2 \rangle \\ & = \frac{1}{(1 - gC(F, G))\sigma_n} \langle \phi_2, (H_n - E_n)(\mathbf{1}^n \otimes P_{\Omega_n})\phi_2 \rangle \\ & \leq \frac{1}{(1 - gC(F, G))\sigma_n} \langle \phi_2, (H_n - E_n)\phi_2 \rangle. \end{aligned}$$

Note that

$$(5.50) \quad E \leq \langle \tilde{\phi}_n, H\tilde{\phi}_n \rangle = \langle \phi_n, H^n\phi_n \rangle = E^n = E_n$$

In view of (5.49) and of (5.50) we get

$$(5.51) \quad \begin{aligned} & \frac{\langle \phi_2, (H_n - E_n)\phi_2 \rangle}{(1 - gC(F, G))\sigma_n} \\ & = \frac{\langle \phi_2, (E - E_n)\phi_2 \rangle}{(1 - gC(F, G))\sigma_n} + \frac{\langle \phi_2, (H_n - H)\phi_2 \rangle}{(1 - gC(F, G))\sigma_n} \\ & \leq \frac{\langle \phi_2, (H_n - H)\phi_2 \rangle}{(1 - gC(F, G))\sigma_n}. \end{aligned}$$

Hence

$$(5.52) \quad \langle \phi_2, E_{\{E_n\}}(H^n)^\perp \otimes P_{\Omega_n}\phi_2 \rangle \leq \frac{\langle \phi_j, (H_n - H)\phi_2 \rangle}{(1 - g_2^{(\delta)}C(F, G))\sigma_n}$$

Here $g_2^{(\delta)}$ has been introduced in proposition 5.3

Estimate of $\|(H_n - H)\tilde{\phi}\|$. We have

$$(5.53) \quad H - H_n = g(H_I - H_{I,n})$$

$H - H_n$ is associated with the kernels $F^{(\beta)}(\xi_2, \xi_3)\chi_{\sigma_n}(\mathbf{p}_4)G^{(\beta)}(\xi_1, \xi_4)$.

By adapting the proof of (5.2) to the estimate of $(H - H_n)$ we finally get

$$(5.54) \quad \|(H - H_n)\phi_2\| = g\|(H_I - H_{I,n})\phi_2\| \leq gK_n(F, G)(C\|H_0\phi_2\| + B)$$

where

$$(5.55) \quad K_n(F, G) = \sum_{\beta=1}^2 \|F^{(\beta)}(\cdot, \cdot)\|_{L^2} \|\chi_{\sigma_n}(\mathbf{p}_4)G^{(\beta)}(\cdot, \cdot)\|_{L^2}.$$

Under Hypothesis 5.2(ii) we get

$$(5.56) \quad K_n(F, G) \leq 2\left(\sum_{\beta=1}^2 \|F^{(\beta)}(\cdot, \cdot)\|_{L^2}\right) \tilde{K}(G)\sigma_n.$$

This, together with (5.55), yields

$$(5.57) \quad |\langle \phi_2, (H - H_n)\phi_2 \rangle| \leq gK\sigma_n.$$

where $K = 2\left(\sum_{\beta=1}^2 \|F^{(\beta)}(\cdot, \cdot)\|_{L^2}\right) \tilde{K}(G)(C\|H_0\phi_2\| + B)$.

Combing (5.41), (5.42), (5.46), (5.52) and (5.57) we finally get

$$(5.58) \quad \langle \phi_2, (\mathbf{1}_{\mathfrak{F}} - E_{\{E_n\}}(H_n))\phi_2 \rangle \leq gK'$$

Here $K' = \frac{K}{(1-g_2^{(\delta)})C(F, G)}$.

K' is a positive constant independent of g and it follows from (5.41) that, for g sufficiently small, $\langle \phi_1, \phi_2 \rangle \neq 0$. This is a contradiction and $p = 1$. This concludes the proof of theorem 5.2. □

ACKNOWLEDGEMENTS.

J.-C. G. acknowledges J.-M Barbaroux, J. Faupin and G. Hachem for helpful discussions.

REFERENCES

- [1] L. Amour, B. Grébert, and J.-C. Guillot. A mathematical model for the Fermi weak interactions. *Cubo*, 9(2),(2007), 37–57.
- [2] A. Arai. Essential spectrum of a self-adjoint opeator on a abstract hilbert of Fock type and applications to quantum field Halmitonians. *J. Math. Anal. Appl.*, 246 ,(2000), 189-216.
- [3] V. Bach, J. Fröhlich, and I. Sigal. Quantum electrodynamics of confined relativistic particles. *Adv. Math.*, 137:299–395, 1998.

- [4] W.H. Aschbacher, J.-M. Barbaroux, J. Faupin and J.-C. Guillot. Spectral theory for a mathematical model of the weak interaction: The decay of the intermediate vector bosons W^\pm . II. *Ann. Henri Poincaré*, 12 , (2011), 1539–1570. [arXiv:1105.2247](#)
- [5] J.-M. Barbaroux, M. Dimassi and J.-C. Guillot. Quantum electrodynamics of relativistic bound states with cutoffs. *J. Hyperbolic Differ. Equ.*, 1(2), (2004), 271–314.
- [6] J.-M. Barbaroux and J.-C. Guillot. Spectral theory for a mathematical model of the weak interaction: The decay of the intermediate vector bosons W^\pm . I. *Advances in Mathematical Physics* ID 978903 (2009), [arXiv:0904.3171](#)
- [7] J.-M. Barbaroux, J. Faupin and J.-C. Guillot. Local decay for weak interactions with massless particles. [arXiv:1611.07814](#)
- [8] J. Beranger *et al* (Particle Data Group). *Phys. Rev.D* , 86 , O10001, 2012.
- [9] K. Bhattacharya. Solution of the Dirac equation in presence of an uniform magnetic field. [arXiv:0705.4275](#)
- [10] K. Bhattacharya and P.B.Pal. Inverse beta decay of arbitrarily polarized neutrons in a magnetic field. *Pramana.*, 62, (2004) , 1041. [arXiv:hep-ph/0209053](#)
- [11] K. Bhattacharya and P.B.Pal. Neutrinos and magnetic fields: a short review. *Proc. Ind. Natl. Sci. Acad.*, 70 , (2004) , 145. [arXiv:hep-ph/0212118](#)
- [12] J. Dereziński and C. Gérard. Asymptotic completeness in quantum field theory. Massive Pauli-Fierz Hamiltonians. *Rev. Math. Phys.*, 11(4) , (1999), 383–450.
- [13] J. Dereziński and C. Gérard. *Mathematics of quantization and quantum fields*. Cambridge University Press. 2013.
- [14] H. Duan and Y.Z. Qian. Rates of neutrino absorption on neutrons and the reverse processes in strong magnetic fields. *Phys. Rev.D72.*, (2005), 023005. [arXiv:astro-ph/0506033](#)
- [15] J. Glimm, A. Jaffe. *Quantum Field Theory and Statistical Mechanics*. Birkhäuser, Boston inc., Boston MA , 1985, Expositions, Reprint of articles published in 1969–1977.
- [16] W. Greiner and B. Müller. *Gauge Theory of weak interactions*. Springer, Berlin, 1989.
- [17] C. Guinti and A. Studenikin. Neutrino electromagnetic interactions: a window to new physics. *Rev. Mod. Phys.*, 87 , 2015 , 531 . [arXiv:1403.6344](#)
- [18] J.C. Guillot. Spectral theory of a mathematical model in quantum field theory for any spin. *Contemporary Mathematics*, 640 , 2015 , 13–37.
- [19] G. Hachem Effet Zeeman pour un électron de Dirac. *Ann. Inst. Henri Poincaré*, 58 , 1993 , 105–123 .
- [20] F. Hiroshima. Multiplicity of ground states in quantum field theory models: applications of asymptotic fields. *J. Funct. Anal.*, 224 , (2005) , 431–470.
- [21] C. Itzykson and J.-B. Zuber. *Quantum Field Theory*. McGraw-Hill book company. 1980.
- [22] M.H. Johnson and B.A. Lippmann. Relativistic motion in a magnetic field. *Phys. Rev.*, 77 ,(1950), 702–705.
- [23] F.A. Reuse. *Electrodynamique et Optiques Quantiques*. Presses Polytechniques et Universitaires Romanes , Lausanne, 2007.
- [24] T. Takaesu. On the spectral analysis of quantum electrodynamics with spatial cutoffs. I. *J. Math. Phys.*, 50 , (2009), 06230.
- [25] T. Takaesu. On generalized spin-boson models with singular perturbations. *Hokkaido Mathematical Journal*, 39 , (2010) , 1–33.
- [26] T. Takaesu. Essential Spectrum of a Fermionic Quantum Field Model. *Infinite Dimensional Analysis, Quantum Probability and Related Topics*, 17 , No4, (2014), 1450024.
- [27] B. Thaller. *The Dirac Equation*. Texts and Monographs in Physics, Springer Verlag, Berlin, 1992.
- [28] S. Weinberg. *The quantum theory of fields. Vol. I*. Cambridge University Press, Cambridge, 2005. Foundations.

[29] S. Weinberg. *The quantum theory of fields. Vol. II.* Cambridge University Press, Cambridge, 2005. Modern applications.

(J.-C. Guillot) CENTRE DE MATHÉMATIQUES APPLIQUÉES, UMR 7641, ÉCOLE POLYTECHNIQUE - C.N.R.S, 91128 PALAISEAU CEDEX, FRANCE
E-mail address: `Jean-Claude.Guillot@polytechnique.edu`