On one-loop impacts of the Rashba coupling

J. R. Nascimento, A. Yu. Petrov\textsuperscript{1} and H. Belich\textsuperscript{2}

\textsuperscript{1}Departamento de Física, Universidade Federal da Paraíba
Caixa Postal 5008, 58051-970, João Pessoa, Paraíba, Brazil\textsuperscript{*}

\textsuperscript{2}Departamento de Física e Química,
Universidade Federal do Espírito Santo, Av. Fernando Ferrari,
514, Goiabeiras, 29060-900, Vitória, ES, Brazil\textsuperscript{†}

Abstract

In this paper, we describe the one-loop contributions in QED with Rashba coupling. We show that all purely nonminimal contributions are explicitly finite, so, the whole theory is one-loop renormalizable.

PACS numbers: 11.30.Cp

\textsuperscript{*}Electronic address: jroberto, petrov@fisica.ufpb.br

\textsuperscript{†}Electronic address: belichjr@gmail.com
I. INTRODUCTION

The electroweak unification performed through the Higgs mechanism was one of the largest achievements of the Standard Model of particle physics (SM). With an idea of spontaneous gauge symmetry violation, the intermediate bosons corresponding to weak interaction gain mass and the photons continues to be massless. To complete the success of experimental predictions, in 2013 the detection of Higgs boson was confirmed. However, it is still necessary to understand why the mass of the Higgs boson is approximately 126 GeV, the problem of the hierarchy of gauge, dark matter, and dark energy. Then we need a theory that goes beyond the Standard Model.

One proposal of investigation is a line of research that deals with the Spontaneous Lorentz Violation (SLV) which is induced by fluctuations of primordial fields in the spacetime which grows up when we aims to arrives at Planck scale ($10^{19}$ GeV). This possibility of SLV allows to introduce at least one privileged direction in the spacetime \cite{1-10}. The first attempt to include SLV in the Standard Model of Particle Physics became known as the Standard Model Extended (SME) \cite{11-13}.

In order to verify the possibility of SLV we have the option to observe a non-trivial background by measurements performed in the particle accelerators and in low-energy scenarios involving quantum mechanical effects \cite{14-19}. Namely the Quantum Mechanical description of a interactive particle with this background can give us a hint for this fundamental theory. By a non-minimal coupling with this environment we can estimate the energy scale at which SLV can emerge through the bounds calculated from the uncertainty of interference experiments \cite{20-24}. This line is based on the idea that these anisotropies can generate new Berry phases \cite{25} acquired by a particle which is moving in this region \cite{26}.

In this paper we follow this line of research describing the spontaneous violation of the Lorentz symmetry caused by a tensor background. Particularly, a non-minimal Lorentz-violating coupling which appears in \cite{27, 28} calls our attention. Using such non-minimal couplings, we discuss the arising of the Landau system and the influence of a Rashba-like coupling induced by a Lorentz symmetry violation scenario in the nonrelativistic quantum dynamics for a spin-1/2 neutral particle. In the ref. \cite{29}, with an analogue of the Landau system for a neutral particle, a interesting bound of the Lorentz breaking term, that is, $g b_0 < 2.2 \times 10^{-6} (eV)^{-3}$ is estimated.
In this paper, we intend to generate nonlinear contributions to the action via an appropriate coupling. In the section 2, we carry out the quantum calculations, and the section 3 is our conclusion.

II. NON-MINIMAL LORENTZ-VIOLATING COUPLING

The Spin-Orbit Coupling (SOC), which is derived from the Dirac equation by the prescription of the Foldy-Wouthuysen Representation \([14–19]\), has recently been revisited in order to understand a number of new proposals for materials with unusual behavior. The emergence of new possibilities of spintronic devices \([30]\) without applying a magnetic field has been the cause nowadays of a intensive research in Condensed Matter \([31]\) and High Energy Physics \([25, 26]\). The SOC can appear in new materials due to 3-dimensional coupling, to the presence of the surface acting as an interface (Rashba and Dresselhaus effect), and due to the 1-dimensional coupling. Our objective in this article is to call attention to new types of coupling in electrodynamics \([27–29]\) and possible scenarios of their manifestation. To be more specific, the problem is whether we can generate such couplings by radiative correction processes. Our objective in this paper is to study the perturbative impacts of Rashba type couplings.

The possibility to go beyond the SM we use consists, by relaxing the renormalizability requirement, in proposing an effective Dirac equation with non-minimal coupling since we are searching for a more fundamental theory \([21, 32, 33]\).

Let us consider the spinor QED with the Rashba coupling \([27]\):

\[
S = \int d^4x \bar{\psi}(i\gamma^\mu \partial_\mu - eA_\mu - \frac{g}{2} F_{\mu\alpha} F^{\alpha}_{\mu\nu} \gamma^{\mu} \gamma^{\nu} - m)\psi.
\]

Although this theory is non-renormalizable, it displays very interesting properties at the one-loop level. Actually, we see that there is no one-loop divergences generated by the non-minimal coupling in this theory. We will find the lower contribution to the one-loop effective action. It is clear that due to the absence of the \(\gamma_5\) matrix and Levi-Civita symbol, we cannot have the Carroll-Field-Jackiw term.

The propagator of the spinor is usual, \(<\psi(k)|\bar{\psi}(-k)> = \frac{1}{k^2 + m}\). Also, we restrict ourselves by constant background electric and magnetic fields, thus imposing the condition \(\partial_\mu F_{\nu\lambda} = 0\).

Then, the contribution with only one vertex, of the first order in \(b_\mu\), evidently gives zero
result. Indeed, one cannot form a scalar from a product of tensors with odd total number
of indices (one index of \( b_\mu \) and even number of indices of any order of \( F_{\mu\nu} \)). Hence, the lower
correction is

\[
\Gamma_4 = \frac{1}{2} (F^2)_{\alpha\gamma} (F^2)_{\mu\rho} b_\beta b_\nu \text{tr} \int \frac{d^4k}{(2\pi)^4} \frac{\gamma^\alpha \gamma^\beta \gamma^\gamma (k^2 + m) \gamma^\mu \gamma^\nu \gamma^\rho (k^2 + m)}{(k^2 - m^2)^2}. \tag{2}
\]

Here \((F^2)_{\mu\nu} = F_{\mu\alpha} F_{\nu\alpha} \). We note the symmetry of \( F^2 \) tensor, i.e.
\((F^2)_{\mu\nu} = (F^2)_{\nu\mu} \). To proceed with the expression \((2)\), we use the formula

\[
\text{tr}(\gamma^\alpha \gamma^\beta \gamma^\gamma \gamma^\delta) = 4(\eta^{\alpha\beta} \eta^{\gamma\delta} - \eta^{\alpha\gamma} \eta^{\beta\delta} + \eta^{\alpha\delta} \eta^{\beta\gamma}), \tag{3}
\]

and go to \( d \)-dimensional space-time, making the replacement \( k_\alpha k_\beta \to \frac{k^2}{d} \eta_{\alpha\beta} \). Also, we keep
only even orders in momenta since the odd ones yield zero contributions. We have for the
numerator \( N_{\alpha\beta\gamma\mu\nu\rho} \):

\[
N_{\alpha\beta\gamma\mu\nu\rho} = \text{tr}[\gamma^\alpha \gamma^\beta \gamma^\gamma (k^2 + m) \gamma^\mu \gamma^\nu \gamma^\rho (k^2 + m)] =
\]

\[
= \gamma^\alpha \gamma^\beta \gamma^\gamma \frac{k^2}{d} \gamma^\mu \gamma^\nu \gamma^\rho k^2 + m^2 \gamma^\alpha \gamma^\beta \gamma^\gamma \gamma^\mu \gamma^\nu \gamma^\rho] + \ldots =
\]

\[
= \text{tr}\left[\frac{k^2}{d} \gamma^\alpha \gamma^\beta \gamma^\gamma \gamma^\mu \gamma^\nu \gamma^\rho \gamma^\sigma + m^2 \gamma^\alpha \gamma^\beta \gamma^\gamma \gamma^\mu \gamma^\nu \gamma^\rho \right] + \ldots, \tag{4}
\]

where dots are for irrelevant terms. Then, we use the relation

\[
\gamma^\sigma \gamma^\mu \gamma^\nu \gamma^\rho = -2 \gamma^\rho \gamma^\nu \gamma^\mu, \tag{5}
\]

which gives

\[
N_{\alpha\beta\gamma\mu\nu\rho} = \text{tr}\left[-\frac{2k^2}{d} \gamma^\alpha \gamma^\beta \gamma^\gamma \gamma^\mu \gamma^\nu \gamma^\rho + m^2 \gamma^\alpha \gamma^\beta \gamma^\gamma \gamma^\mu \gamma^\nu \gamma^\rho \right] + \ldots. \tag{6}
\]

We remind that this trace should be contracted with \((F^2)_{\alpha\gamma} (F^2)_{\mu\rho} b_\beta b_\nu \), and \((F^2)_{\mu\rho} \) is symmetric. Actually, it means that the equivalent trace for \( N_{\alpha\beta\gamma\mu\nu\rho} \), under replacement of \( \mu \) by \( \rho \) in one of the terms, is

\[
N_{\alpha\beta\gamma\mu\nu\rho} \simeq (-\frac{2k^2}{d} + m^2) \text{tr}[\gamma^\alpha \gamma^\beta \gamma^\gamma \gamma^\mu \gamma^\nu \gamma^\rho] + \ldots. \tag{7}
\]

Then, we consider \( \Gamma_4 \), which after these replacements takes the form

\[
\Gamma_4 = \frac{1}{2} (F^2)_{\alpha\gamma} (F^2)_{\mu\rho} b_\beta b_\nu \text{tr} \int \frac{d^4k}{(2\pi)^4} \frac{(-\frac{2k^2}{d} + m^2) \text{tr}[\gamma^\alpha \gamma^\beta \gamma^\gamma \gamma^\mu \gamma^\nu \gamma^\rho]}{(k^2 - m^2)^2}. \tag{8}
\]

Then, it is easy to see that

\[
(F^2)_{\alpha\gamma} b_\beta \gamma^\alpha \gamma^\beta \gamma^\gamma = 2 (F^2)_{\alpha\gamma} b_\gamma - \partial (F^2)_{\alpha}. \tag{9}
\]
So, we have

\[ \Gamma_4 = \frac{1}{2} \int \frac{d^4k}{(2\pi)^4} \frac{(-k^2 + m^2)^2}{(k^2 - m^2)^2} \text{tr}\{[2(F_2)^\alpha \gamma^\beta b^\alpha \gamma^\beta - \beta(F_2)^\alpha_{\alpha}\}[2(F_2)^\beta b^\beta \gamma^\delta - \beta(F_2)^\beta_{\beta}]\}. \]  

(10)

Calculating the trace, we find

\[ \Gamma_4 = \frac{1}{2} \int \frac{d^4k}{(2\pi)^4} \frac{(-2k^2 + m^2)^2}{(k^2 - m^2)^2} [16(F_2)^\alpha \gamma^\beta b_\alpha(F_2)^\gamma_{\beta\gamma} - 4b^2(F_2)^\beta_{\beta}(F_2)^\alpha_{\alpha} - 16(F_2)^2_{\alpha\beta} b^\alpha b^\beta(F_2)^\gamma_{\gamma}]. \]  

(11)

It remains to integrate over the momenta. We do the Wick rotation:

\[ I = \frac{1}{2} \int \frac{d^dk}{(2\pi)^d} \frac{(-2k^2 + m^2)^2}{(k^2 - m^2)^2} = \frac{i}{d} \int \frac{d^dk_E}{(2\pi)^d} \frac{k^2 + \frac{d}{2}m^2}{(k^2_E + m^2)^2}. \]  

(12)

However, this is nothing more than the miraculous integral from \[32\] which is finite despite naively it involves both quadratic and logarithmic divergences. For the first manner of calculating this integral, we choose \( d = 4 \). In this case, similarly to \[32\], we find that

\[ I = \frac{im^2}{64\pi^2}. \]  

(13)

Within this prescription, we get after the inverse Wick rotation

\[ \Gamma_4 = \frac{m^2}{64\pi^2}[16(F_2)^\alpha \gamma^\beta b_\alpha(F_2)^\gamma_{\beta\gamma} - 4b^2(F_2)^\beta_{\beta}(F_2)^\alpha_{\alpha} - 16(F_2)^2_{\alpha\beta} b^\alpha b^\beta(F_2)^\gamma_{\gamma}]. \]  

(14)

We note that this integral is ambiguous, so, other prescriptions for its calculation yield other results.

Within the second manner, we choose \( d = 4 - \epsilon \). Afterward, our integral \( I \) is

\[ I = \frac{i}{4 - \epsilon} \int \frac{d^{4-\epsilon}k_E}{(2\pi)^{4-\epsilon}} \frac{k^2 + \frac{4-\epsilon}{2}m^2}{(k^2_E + m^2)^2}. \]  

(15)

For a small \( \epsilon \neq 0 \) this integral is equal to zero, cf. \[33\]. Hence, our four-point function is ambiguous.

The next correction to be studied involves four vertices (the contributions with three vertices will vanish identically unless we consider derivatives of external fields). However, it is equivalent to the four-leg diagram of the usual QED where the replacement of the external legs by the rule \( A_\gamma \rightarrow 2(F_2)^\alpha \gamma^\beta \rightarrow \gamma^\beta b^\gamma - b_\gamma(F_2)^\alpha_{\alpha} \) is carried out. And the four-leg diagram in QED (with no derivatives on external fields) is well known to vanish.
We note that the divergent contributions from Feynman diagrams with three (similarly, five, seven etc.) ”new” vertices will also vanish – indeed, it is well known that the one-loop three-point function of the gauge field in QED yields zero result, and also, we can remind that an $n$-point in our theory can be obtained from that one in QED through the replacement $A_\gamma \rightarrow 2(F^2)_{\alpha\gamma}b^\alpha - b_\gamma(F^2)^\alpha_\alpha$, while the integral over momenta is just the same as in QED. In other words, since the one-loop three-point function in QED vanishes in some regularization to provide the gauge symmetry, we can conclude that all divergent contributions from diagrams with three vertices in our theory will vanish as well.

The (one-loop) graphs with six and more ”new” vertices are explicitly finite. Actually, we showed that the theory \( \Pi \) is one-loop finite at $e = 0$. We notice that in [27], namely the case $e = 0$ was treated. At the same time, it is interesting to discuss the one-loop behavior of the complete theory \( \Pi \) involving both $e$ and $g$ which is an extension of the QED.

First of all, it is clear that the total number of vertices in any one-loop Feynman diagram with external gauge legs must be even (indeed, each vertex carries an odd number of indices of fields and $b_\mu$ vectors contracted to some Dirac matrices), and the number of the external $bF^2$ lines should be even as well, otherwise we should have an odd number of $A_\alpha$ legs without derivatives, which is inconsistent with the gauge invariance. Second, since our vertices do not involve derivatives acting on spinor fields, the upper limit for the degree of divergence of an arbitrary one-loop graph is $\omega = 4 - V$, where $V$ is a number of vertices. Actually, it is even less if some derivatives are transported to external fields, being $\omega = 4 - V - 2N_d$, where $N_d$ is the number of derivatives which do not present in vertices from the beginning but arise within the derivative expansion. It means that the only one-loop superficially divergent graphs in the theory are those ones with $V = 2$ or $V = 4$, and, moreover, a divergent graph with four vertices cannot contain any minimal vertex because, in order to form the gauge invariant combination, some derivatives must be moved to external fields associated with these vertices which decreases the degree of divergence of the corresponding Feynman diagram (for example, it is well known that the logarithmically divergent contribution from the $A^4$ graph with four minimal vertices is not gauge invariant and hence vanishes, and, replacing the external legs by the rule $A_\gamma \rightarrow 2(F^2)_{\alpha\gamma}b^\alpha - b_\gamma(F^2)^\alpha_\alpha$, we find that the same situation occurs for $(bF^2)^4$ contribution), and thus the corresponding Feynman diagram becomes finite. Moreover, there is no way to generate the CFJ term since there is neither $\gamma_5$ matrices nor Levi-Civita symbol in the classical Lagrangian of the theory, so, the quantum
corrections should be some contractions of $F_{\mu\nu}$ tensors only.

Therefore we see that actually, at the one-loop order we can have only the divergent corrections proportional to $(bF^2)^2$ or $(bF^2)^4$, besides of the usual $F^2$. At the same time, we note that the $(bF^2)^4$ contribution vanishes under some regularization prescription, as we already argued, and the $(bF^2)^2$ contribution was showed above to be explicitly finite. So, we see that actually the only one-loop divergence in this theory is just that one occurring in the usual QED (see f.e. [34]), and there is no new one-loop divergences generated by minimal vertices. Actually, we showed that the theory involving both usual and Rashba-like couplings is one-loop renormalizable in the gauge sector, with the only divergence is that one arising in an usual QED.

III. SUMMARY

We considered the Lorentz-breaking extended QED with the additional Rashba coupling. It turns out to be that in the one-loop approximation, this coupling does not generate any new divergences in the gauge sector, other than that one arising in the usual QED, therefore the resulting theory is one-loop renormalizable, and, for $e = 0$, even one-loop finite. This allows to treat this coupling as an important ingredient of a possible Lorentz-breaking extended QED, at least if we disregard higher-loop corrections, or treat the gauge field as purely external one, thus restricting ourselves by the fermionic determinant (nevertheless, we note that non-renormalizable field theory models are intensively used even outside of the fermionic determinant context, see f.e. [35]). Moreover, this theory allows to generate the nonpolynomial effective action of the gauge field (that is, the Euler-Heisenberg action), with all terms, at $e = 0$, will be explicitly finite.

We note that all studies we carried out here can be applied as well if we replace $F_{\mu\nu}$ by its dual $\tilde{F}_{\mu\nu}$ in the fermion-vector coupling, as it has been done in [29] within the geometrical phase context. So, the theory with the coupling on the base of $\tilde{F}$ displays the same one-loop properties as the theory we considered in the paper. In principle, other non-minimal spinor-vector couplings can be also studied in this manner.

Acknowledgements. This work was partially supported by Conselho Nacional de Desenvolvimento Científico e Tecnológico (CNPq). The work by A. Yu. P. has been supported
by the CNPq project No. 303783/2015-0.

[1] H. Belich, T. Costa-Soares, M. A. Santos and M. T. D. Orlando, Rev. Bras. Ensino Fís. 29, 1 (2007).
[2] A. P. Baeta Scarpelli, H. Belich, J. L. Bokdo, L. P. Colatto, J. A. Helayël-Neto, A. L. M. A. Nogueira, Nucl. Phys. Proc. Suppl. 127, 105 (2004).
[3] V. W. Hughes et al, Phys. Rev. Lett. 87, 111804 (2001); R. Bluhm, V. A. Kostelecký, and C. D. Lane, Phys. Rev. Lett. 84, 1098 (2000); E. O. Iltan, JHEP 0306, 016 (2003).
[4] Y. B. Hsiung et al., Nucl. Phys. Proc. Suppl. 86, 312 (2000); FOCUS Collaboration, J. M. Link et al, Phys. Lett. B 556, 7 (2003).
[5] V. A. Kostelecký, Phys. Rev. D 64, 076001 (2001); V. A. Kostelecký and R. J. Van Kooten, Phys. Rev. D 82, 101702(R) (2010).
[6] F. Allmendinger et al, Phys. Rev. Lett. 112, 110801 (2014); B. M. Roberts, Y. V. Stadnik, V. A. Dzuba, V. V. Flambaum, N. Leefer and D. Budker, Phys. Rev. D 90, 096005 (2014); B. M. Roberts, Y. V. Stadnik, V. A. Dzuba, V. V. Flambaum, N. Leefer and D. Budker, Phys. Rev. Lett. 113, 081601 (2014).
[7] V. A. Kostelecký and M. Mewes, Phys. Rev. Lett. 110, 201601 (2013).
[8] V. A. Kostelecký and M. Mewes, Phys. Rev. Lett. 87, 251304 (2001); V. A. Kostelecký and M. Mewes, Phys. Rev. D 66, 056005 (2002).
[9] R. Lehnert, J. Math. Phys. 45, 3399 (2004); B. Altschul, Phys. Rev. Lett. 96, 201101 (2006).
[10] J. S. Díaz and V. A. Kostelecký, Phys. Rev. D 85, 016013 (2012); J. S. Díaz, V. A. Kostelecký, and R. Lehnert, Phys. Rev. D 88, 071902(R) (2013); J. S. Díaz, V.A. Kostelecký and M. Mewes, Phys. Rev. D 89, 043005 (2014).
[11] V. A. Kostelecký and S. Samuel, Phys. Rev. D 39, 683 (1989).
[12] D. Colladay and V. A. Kostelecký, Phys. Rev. D 55, 6760 (1997).
[13] D. Colladay and V.A. Kostelecký, Phys. Rev. D 58, 116002 (1998).
[14] H. Belich and K. Bakke, Phys. Rev. D 90, 025026 (2014).
[15] V. A. Kostelecký and N. Russell, Rev. Mod. Phys. 83, 11 (2011).
[16] M. M. Ferreira Jr. and F. M. O. Moucherek, Nucl. Phys. A 790, 635c (2007).
[17] K. Bakke and H. Belich, Ann. Phys. (Berlin) 526, 187 (2013).
[18] H. Belich, F. J. L. Leal, H. L. C. Louzada and M. T. D. Orlando, Phys. Rev. D 86, 125037 (2012).

[19] R. Casana, M. M. Ferreira, Jr., E. Passos, F. E. P. dos Santos and E. O. Silva, Phys. Rev. D 87, 047701 (2013); J. B. Araujo, R. Casana, M. M. Ferreira, Phys. Lett. B 760, 302 (2016); R. Casana, M. M. Ferreira, A. L. Mota, Ann. Phys. 375, 179 (2016); R. Casana, M.M. Ferreira, F. E. P. dos Santos, Phys. Rev. D 94, 125011 (2016).

[20] H. Belich, T. Costa-Soares, M. M. Ferreira Jr. and J. A. Helayël-Neto, Eur. Phys. J. C 41, 421 (2005).

[21] H. Belich, T. Costa-Soares, M. M. Ferreira Jr., J. A. Helayël-Neto, M. T. D. Orlando, Phys. Lett. B 639, 675 (2006).

[22] K. Bakke and H. Belich, J. Phys. G: Nucl. Part. Phys. 40, 065002 (2013).

[23] K. Bakke and H. Belich, J. Phys. G: Nucl. Part. Phys. 39, 085001 (2012).

[24] A. G. de Lima, H. Belich and K. Bakke, Eur. Phys. J. Plus 128, 154 (2013).

[25] K. Bakke and C. Furtado, Phys. Rev. D 80, 024033 (2009); K. Bakke and C. Furtado, Phys. Lett. A 375, 3956 (2011); K. Bakke and C. Furtado, Quantum Inf. Comput. 11, 444 (2011); K. Bakke, A. Yu. Petrov, C. Furtado, Annals Phys. 327, 2946 (2012); K. Bakke and C. Furtado, Ann. Phys. (NY) 327, 376 (2012); K. Bakke and C. Furtado, Quantum Inf. Process. 12, 119 (2013); H. F. Mota, K. Bakke, Phys. Rev. D 89, 027702 (2014).

[26] K. Bakke, H. Belich, J. Phys. G, 40, 065002, (2013); K. Bakke and H. Belich, Annalen der Physik (Leipzig) 526, 187, (2014); A. G. de Lima, H. Belich and K. Bakke, Annalen der Physik (Leipzig), 526, 514 (2014); H. Belich, K. Bakke, Phys. Rev. D, 90, 112 (2014); K. Bakke, H. Belich, Eur. Phys. J. Plus, 129, 147 (2014); H. Belich, K. Bakke, Int. J. Mod. Phys. A 30, 1550136 (2015); K. Bakke, H. Belich, J. Phys. G, 42, 095001 (2015); K. Bakke, H. Belich, Ann. Phys. (NY), 360, 596 (2015), H. Belich, M. A. Santos, M. T. D. Orlando, Mod. Phys. Lett. A 30, 1550191 (2015).

[27] K. Bakke, H. Belich, J. Phys. G 39, 085001 (2012).

[28] K. Bakke, H. Belich, Ann. Phys. (NY), 333, 272 (2013).

[29] K. Bakke, H. Belich, Ann. Phys. (NY), 354, 1 (2015).

[30] P. D. Johnson, Rep. Prog. Phys. 60, 1217 (1997).

[31] Y. A. Bychkov, E. I. Rashba, J. Phys. C 17, 6039 (1984).

[32] M. Gomes, J. R. Nascimento, A. Yu. Petrov, A. J. da Silva, Phys. Rev. D 81, 045018 (2010).
[33] T. Mariz, J. R. Nascimento, A. Yu. Petrov, A. P. Baeta Scarpelli, Eur. Phys. J. C 73, 2526 (2013).

[34] K. Fujikawa, Path Integrals and Quantum Anomalies, Clarendon Press, Oxford, 2004.

[35] A. Brignole, Nucl. Phys. B 579, 101 (2000).