NEW INTERPRETATION OF EXPERIMENTS ON THE INTERMEDIATE MESON

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A new theory of weak interactions is proposed in which the coupling between the $V-A$ currents $j_\mu(X)$ and $j'_\mu(X')$ is achieved not by vector mesons [by a propagator $D_{\mu\nu}^{\pm}(X - X')$], but by a scalar function $R(X - X')$, a fermion—antifermion loop which plays the role of a unique film joining two different points $X$ and $X'$ of completely uncoupled space—times (as a result of which the space becomes a continuum). The existence of the actual currents $j_\mu$ results from correlations between the two different spinor layers of Dirac layer formation.

INTRODUCTION

It is well known that the original theory of weak interactions proposed by Fermi was based on local couplings of four fermion fields $\psi(X)$ (see Fig. 1a). With further refinement, resulting from taking account of $P$- and $C$-asymmetry, there remained only couplings of the type $j_\mu(X)j'_\mu(X')$ where $j_\mu = \frac{1}{2}O_\mu\psi$ is the $V-A$ current [of the matrix $O_\mu = i\gamma_\mu(1 + \gamma_5)$]. In recent times, in view of the adoption of the hypothesis (it must be said, poorly based physically) of a unified nature of weak and electromagnetic interactions, the principal coupling responsible for weak interactions has been written (as also has electromagnetic coupling) in the form $j_\mu W_\mu$, where $W_\mu$ are the fields of the so-called intermediate vector mesons (these are the $W^{\pm}$ and $Z^0$ mesons and also photons). Consequently, for Fermi coupling one obtains a nonlocal expression of the form $\cdot C j_\mu(X)D_{\mu\nu}(X - X')j_\nu(X')$ (Fig. 1b), where $D_{\mu\nu}$ is the vector field propagator.

A weak spot in the intermediate meson theory is that it is based on the concept of the existence of so-called Higgs mesons which are defined by many parameters in the theory (for example, the masses of the $W$ and $Z$ mesons and their very existence, and also the masses of the leptons) which, however, have not yet been observed experimentally. This situation seriously undermines confidence in the theory and its adequacy to the very nature of things (if a Higgs vacuum does not exist, then neither does the intermediate meson). In this case the urgent necessity arises of reinterpreting the results of known experiments which indicate that maxima exist in the collision cross section in the region of $\sqrt{\kappa^2} = 70-90$ GeV [1] (where $\sqrt{\kappa^2}$ is the total energy of the particles in the opposing beams) [1].

An attempt is made in the present paper to undertake such a reinterpretation. We shall first show that the mathematical apparatus of the proposed theory makes it possible to predict the existence of the above-mentioned maxima without invoking a hypothesis concerning the existence of intermediate vector mesons. It turns out that the pole approximation used here is fully justified only for $\kappa^2 < m^2$ (where $m$ is the fermion mass in the fermion—antifermion loop, see Fig. 1c). For $\kappa^2 > m^2$ the behavior of the cross section is entirely different (see Fig. 2), as a function of $\kappa^2$ the cross section is not symmetric about the point $\kappa^2 = m^2$.

We shall preface the formulation of the basic postulates of the proposed theory (Sec. 2) with one mathematical preparation, evidently clear in itself without any theory.
I. MATHEMATICAL PREPARATION

Here we perform a calculation of the fermion–antifermion loop (Fig. 1c) in the theory of bilocal fields in the case of very heavy virtual particles, \( m \gg 1 \) (the corresponding calculations in the case of light virtual particles were performed in [2]). Such particles, even if they exist, will do so only at very small distances where no types of motion (displacements, space–time evolution) are possible. There they are somehow described by one two-component spinor, let us say \( \varphi \); the other spinor \( \chi \) obtained from \( \varphi \) in accordance with the Dirac equations by differentiation \( \sigma_\mu (\partial/\partial X_\nu) \varphi = m \chi \), has not yet appeared (there is no evolution). In this case, the very same projection operator, for example \( P_+ = (1/2)(1 + \gamma_3) \), will be located at each of the vertices and a vertex itself can only be an even combination of the Dirac matrices \( 1, \gamma_3, \gamma_5 \) (odd combinations of the matrices are not allowed because there is no evolution). All such vertices produce the same result, and we shall therefore consider the simplest case 1. In this case the loop is described by the expression

\[
\Pi (\kappa^2) = iC \int \frac{\mathrm{Sp} \, P_+ (p + m) \, P_+ (p + \kappa + m)}{(p^2 - m^2)((p + \kappa)^2 - m^2)} \, \rho(p^4) \, d^4 p =
\]

(1)

(the kinetic part does not operate), where the form factor is \( \rho(p^2) = (\sin p^2/p^2) \) (see [2]). In [2], in the case of small \( m^2 \), the representation \( \sin p^2 = \frac{1}{2} \int I e^{ip^2 \, da} = \frac{1}{2} \int e^{ip^2 \, da} \), was used for \( \rho(p^2) \) where the interval is \( I = [-1, 1] \). However, in the case of large \( m^2 \) one must use another representation

\[
\sin \frac{p^2}{p^2} = \frac{1}{2} \int \frac{e^{ip^2 \, da}}{C_1} = \frac{1}{2} \left( \int \frac{1}{-1} e^{iap^2 \, da} \right),
\]

(2)

where \( C_1 \) is the complement of \( I \) in \( R \): \( I \cup C_1 = R \). The two representations are almost equivalent since the integral \( \frac{1}{2} \int e^{ia} \, da = \pi \delta(p^2) \) is almost equal to zero [the magnitude of \( \Pi \) calculated from \( \delta(p^2) \) is equal to zero]. The normalization factor in Eq. (1) is chosen so that \( \Pi(0) = 1 \).

Using the representation (2), and also the Fock representation of the propagator

\[
\frac{1}{p^2 - m^2} = -\frac{i}{2} \int_0^\infty e^{ip^2 \, da} = \frac{1}{2} \int_0^\infty e^{iap^2 \, da},
\]

and performing the integration with respect to \( p \) in Eq. (1), we arrive at the expression

\[
\Pi (\kappa^2) = -C_2 m^2 \pi^2 \int_{C_1} \frac{da}{2} \int_0^\infty \frac{\, da \, \bar{b}}{(a + \bar{b} + a)^2} \, e^{-i(\kappa z + \kappa \bar{z}) + i(\kappa \bar{z} - \kappa z)}.
\]

After integration with respect to \( a \) we obtain

\[
\Pi (\kappa^2) = -C_2 m^2 \pi^2 \int_0^1 \frac{dz}{2} \int_0^\infty \frac{\, da}{(a + 1)^2} \, \frac{\sin \alpha \left( m^2 - \kappa^2 \bar{z} + \kappa^2 \frac{\alpha z}{\alpha + 1} \right)}{\alpha \left( m^2 - \kappa^2 z + \kappa^2 \frac{\alpha z}{\alpha + 1} \right)}.
\]

where we have converted from the integration variables \( \alpha \) and \( \beta \) to the new variables \( \sigma \) and \( z \): \( \alpha = \sigma(1 - z), \beta = \sigma z \). Since, for large \( m^2 \gg 1 \), the main contribution to the integral with respect to \( \sigma \) comes from the region of small \( \sigma \ll 1 \), the term \( \kappa^2 \sigma z^2/(\sigma + 1) \) can be neglected. In this case we obtain