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INTERNATIONAL CENTRE FOR
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T.W.B. Kibble

Abdus Salam

and

J. Strathdee

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T.W.B. Kibble
Imperial College, London, England,

Abdus Salam
International Centre for Theoretical Physics, Trieste, Italy
and
Imperial College, London, England,

and

J. Strathdee
International Centre for Theoretical Physics, Trieste, Italy.

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ABSTRACT

We analyse the effects of the intensity-dependent mass shift predicted for electrons in an electromagnetic plane wave on bound states. It is shown that the null result of the experiment of Mowat et al. on the ^{133}Cs ground-state hyperfine splitting was to be expected, although their hypothetical explanation of it is incorrect. The recent suggestion that an intense laser beam might effect the restoration of a spontaneously broken symmetry is examined in more detail, and it is demonstrated that although a pure plane will will not suffice, a superposition of plane waves, for example a standing wave, would do so. At present, however, the effect is unobservably small.

I. INTRODUCTION

Some years ago a mass shift of free electrons in an intense plane-wave electromagnetic field was predicted [1)-3)] and became the subject of lengthy controversy (summarized with full references in Ref.4). Later it was suggested [5)] that a similar effect might exist for bound electrons, leading to intensity-dependent shifts in the frequencies of spectral lines. However, an experimental search for such a shift in the hyperfine splitting of the ^{133}Cs ground state yielded negative results [6)]. On a naive classical argument, this negative conclusion was to be expected, because the mass shift of a free electron can be understood [7)] as a mean kinetic energy of the large amplitude oscillation induced by the field, which is precluded for bound electrons. The quantum-mechanical explanation, on the other hand, is less clear. Mowat et al. [6)] suggested that the e^2A^2 term which is responsible should be regarded as a contribution to the energy rather than the mass, so yielding an unobservable uniform shift of all levels. However, this seems a rather unnatural interpretation in view of the fact that in the Klein-Gordon (or second-order Dirac) equation the e^2A^2 term is clearly an addition to m^2. Tryon [8)] in a treatment of quantum electrodynamics in a coherent photon background finds lowest-order energy level shifts independent of the mass shift, but neither the underlying reasons nor the extent of validity of this result are wholly clear.

In this paper we re-examine the status of the electron mass shift and show that, while the e^2A^2 term alone would yield level shifts corresponding to a change in mass rather than a uniform addition, the effect is almost exactly cancelled in the long wavelength limit by the second-order shift arising from the ep-A term.

Recently [9)] two of us suggested a different role for the mass shift, as a possible mechanism for the restoration of spontaneously broken symmetries, as we showed in the case of intense magnetic fields. Here we examine this possibility in more detail.

The mass shift directly affects only charged particles, whereas what is relevant in deciding whether a symmetry is broken is the effective potential evaluated as a function of the vacuum expectation value of some neutral field. An effect can be produced by virtual charged-particle loops but since the neutral particle is thus to be treated as a bound state, one might anticipate that the effect would be suppressed as in the case of an electron bound in an atom.

There is a simple invariance argument which shows that in a true plane wave the suppression is in fact complete. By gauge invariance, the effective potential $V(\phi,A_\mu)$ must depend on the plane-wave vector potential A_μ only through the field tensor $F_{\mu\nu}$. But it must also be a scalar, and for a plane wave no scalar can be formed from $F_{\mu\nu}$, so that V must be a function of the vacuum expectation value ϕ alone. This conclusion will be verified by more explicit examination of the effective potential. It follows that a spontaneously broken symmetry cannot be restored by a simple plane-wave laser beam, as has also been pointed out by Becker [10)]. However the argument clearly fails for other configurations. In particular, a contribution to the effective potential can be expected from a superposition of two plane waves in different directions, for example from a standing wave produced by reflecting a laser beam from a mirror, since the invariant $F^{\mu\nu}F_{\mu\nu}$ is non-zero in that case.

It may also be worth remarking that, even in a pure plane-wave beam, symmetries might be broken (not restored) by a different mechanism, in which a symmetric tensor field acquires a vacuum expectation value proportional to $F_\mu^{\rho}F_{\rho\nu}$, or equivalently $A^2k_\mu k_\nu$.

It has been pointed out by Neville and Rohrlich [11)] that a null-plane formalism is particularly well suited to the treatment of processes in a plane-wave field. In Sec.II we introduce a version of this formalism closely based on the work of Bjorken, Kogut and Soper [12)], in which the relativistic treatment closely parallels the non-relativistic. However we use a modified (non-orthogonal) co-ordinate system recently employed by Bell and Ruegg [13)] in a treatment of the hydrogen atom. This method is interesting in its own right and may well have other applications.

In Sec.III we show that the electron Green's function in a plane-wave field has a particularly simple form in these co-ordinates, and use it to examine the contribution to an effective potential function. The conclusions are summarized and discussed in Sec.IV.

II. DIRAC EQUATION IN NULL-TIME CO-ORDINATES

We use the system of co-ordinates

$$\tau = t - z, \quad \zeta = z, \quad \underset{\sim}{x} = (x,y), \tag{1}$$

comprising a null-time co-ordinate together with the three cartesian spatial co-ordinates. The distinction between ζ and z is required when we come to the contragredient transformations exemplified by the partial derivatives:

$$\partial_\tau = \partial_t , \quad \partial_\zeta = \partial_z + \partial_t , \quad \underset{\sim}{\partial} = (\partial_x, \partial_y) \ .$$

(In using the two-vector notation we regard all vectors other than $\underset{\sim}{\partial}$ as contravariant. Thus $p_\mu = i\partial_\mu$ but $\underset{\sim}{p} = -i\underset{\sim}{\partial}$.) The metric tensor and its increase in the co-ordinate system (1) are:

$$g_{\mu\nu} = \begin{pmatrix} 1 & 1 & \cdot & \cdot \\ 1 & \cdot & \cdot & \cdot \\ \cdot & \cdot & -1 & \cdot \\ \cdot & \cdot & \cdot & -1 \end{pmatrix} , \quad g^{\mu\nu} = \begin{pmatrix} \cdot & 1 & \cdot & \cdot \\ 1 & -1 & \cdot & \cdot \\ \cdot & \cdot & -1 & \cdot \\ \cdot & \cdot & \cdot & -1 \end{pmatrix} .$$

Since γ^τ is singular, two of the four components of the Dirac equation take the form of constraints so that it can be reduced [12] to a two-component form which in our co-ordinates reads:

$$2\,\pi_\tau\,\psi(x) = \left[\pi_\zeta + (m - i\underset{\sim}{\sigma}\cdot\underset{\sim}{\pi})\,\pi_\zeta^{-1}\,(m + i\underset{\sim}{\sigma}\cdot\underset{\sim}{\pi}) \right] \psi(x) . \tag{2}$$

Here

$$\pi_\mu = p_\mu - eA_\mu(x) = i\partial_\mu - eA_\mu(x) \tag{3}$$

and π_ζ^{-1} is the integral operator defined by

$$\pi_\zeta^{-1} f(x) = \frac{1}{2i} \int d\zeta' \, \varepsilon(\zeta - \zeta') \exp\left[-ie \int_{\zeta'}^{\zeta} d\zeta'' A_\zeta(\tau, \zeta'', \underset{\sim}{x}) \right] f(\tau, \zeta', \underset{\sim}{x}) .$$

Let us first suppose that the external field is time-independent, $A_\mu = A_\mu(\zeta, \underset{\sim}{x})$. Then Eq.(2) has stationary solutions of the form

$$\psi(x) = u_n(\zeta, \underset{\sim}{x}) \, e^{-iE_n \tau} .$$

The u_n are eigenfunctions of the Hamiltonian

$$H = eA_\tau + \tfrac{1}{2}\pi_\zeta + \tfrac{1}{2}(m - i\underset{\sim}{\sigma}\cdot\underset{\sim}{\pi})\,\pi_\zeta^{-1}\,(m + i\underset{\sim}{\sigma}\cdot\underset{\sim}{\pi}) ,$$

and are obtainable from the familiar Dirac wave functions by multiplying by $e^{-iE_n\zeta}$ and applying the projector $P_+ = \frac{1}{2}\gamma^\zeta\gamma^\tau$. The bound-state wave functions may be normalized so that

$$\int d\zeta \, d^2\underset{\sim}{x} \, u_n^\dagger(\zeta, \underset{\sim}{x}) \, u_{n'}(\zeta, \underset{\sim}{x}) = \delta_{nn'} . \tag{4}$$

The scattering states are labelled by the three-momentum p and z component of spin s . We use η to denote p_ζ , so that $p = (\eta, \underset{\sim}{p})$. The covariant normalization of these states is

$$\int d\zeta \, d^2\underset{\sim}{x} \; u_{p's'}^\dagger(\zeta, \underset{\sim}{x})\, u_{ps}(\zeta, \underset{\sim}{x}) = 2|\eta|\,(2\pi)^3\,\delta(\eta' - \eta)\,\delta_2(\underset{\sim}{p}' - \underset{\sim}{p})\,\delta_{s's} . \tag{5}$$

For a free particle,

$$u_{ps}(\zeta, \underset{\sim}{x}) = |2\eta|^{1/2} \, e^{i(\underset{\sim}{p}\cdot\underset{\sim}{x} - \eta\zeta)} \, w_s ,$$

where w_s is a two-component spinor, $\begin{pmatrix}1\\0\end{pmatrix}$ or $\begin{pmatrix}0\\1\end{pmatrix}$. The corresponding energy eigenvalue is

$$E_p = \frac{1}{2}\eta + (m^2 + \underset{\sim}{p}^2)/2\eta . \tag{6}$$

For $\eta < 0$, u_{ps} is of course a negative-energy eigenfunction.

Now let us consider in addition to the time-independent field A_μ a plane wave field $a_\mu(\tau)$. Using standard perturbation methods, we look for a solution to the Dirac equation of the form:

$$\varphi(x) = \sum_n c_n(\tau)\, u_n(\zeta, \mathbf{x}) \quad .$$

Choosing a gauge in which only the transverse components $\mathbf{a}$ of a_μ are non-zero, we find the equation

$$i\,\frac{dc_n(\tau)}{d\tau} = E_n\, c_n(\tau) + e\,\mathbf{a}(\tau)\cdot \sum_{n'} \mathbf{V}_{nn'}\, c_{n'}(\tau) + e^2 \mathbf{a}^2(\tau) \sum_{n'} W_{nn'}\, c_{n'}(\tau), \tag{7}$$

which coincides in form with its non-relativistic counterpart, save for the definitions of $\mathbf{V}$ and W. These are:

$$W_{nn'} = \tfrac{1}{2} \int d\zeta\, d^2\mathbf{x}\; u_n^\dagger(\zeta,\mathbf{x})\, \pi_\zeta^{-1}\, u_{n'}(\zeta,\mathbf{x}) \tag{8}$$

and

$$\mathbf{V}_{nn'} = -\tfrac{1}{2} \int d\zeta\, d^2\mathbf{x}\; u_n^\dagger(\zeta,\mathbf{x}) \Big(\{ \pi_\zeta^{-1}, \boldsymbol{\pi} \} + i e \sigma_3 \left[\pi_\zeta^{-1}, {}^*\boldsymbol{\pi} \right] \Big)\, u_{n'}(\zeta,\mathbf{x}) \; , \tag{9}$$

where π_μ is still defined by (3) and ${}^*\boldsymbol{\pi}$ denotes the dual two-vector, ${}^*\pi^k = \epsilon^{k\ell}\, \pi^\ell$.

For a free particle both $\mathbf{V}$ and W are diagonal in the momentum representation. In fact,

$$\mathbf{V} = -\mathbf{p}/\eta \quad , \quad W = 1/2\eta \; .$$

Hence we recover the solutions found earlier [14),1),3),15)] but in a simpler two-component form:

$$c_{ps}(\tau) = c_{ps}(0) \exp\left(-i \int_0^\tau d\tau'\; \hat{E}_p(\tau') \right) ,$$

with

$$\hat{E}_p(\tau) = E_p - \frac{e\,\mathbf{p}\cdot\mathbf{a}(\tau)}{\eta} + \frac{e^2 \mathbf{a}^2(\tau)}{2\eta}$$
$$= \tfrac{1}{2}\eta + \frac{[\mathbf{p} - e\,\mathbf{a}(\tau)]^2 + m^2}{2\eta} \; . \tag{10}$$

The absence of any explicit spin-dependent factor, similar to the one that appears in the four-component solution, is particularly interesting.

There is one crucial difference between (7) and the corresponding non-relativistic formula, namely that W as given by (8) is no longer a constant $(1/2m)$ but the matrix element of an operator $(1/2\pi_\zeta)$. Thus it yields an energy shift which is not constant but given approximately by

$$\Delta E_n = e^2 \langle \mathbf{a}^2 \rangle_{av}\; W_{nn} \approx e^2 \langle \mathbf{a}^2 \rangle_{av}\; \frac{E_n}{2m^2} \quad , \tag{11}$$

just the shift one would expect from a change in the electron mass.

If the frequency ω of the plane wave is small compared with a characteristic atomic frequency ω_0, then one may expect higher-order terms in the perturbation series to be negligible because of their large energy denominators. However, (11) cannot represent the true level shift, since for a wave of fixed intensity it tends to infinity in the long-wavelength limit. In fact, the matrix elements of $\mathbf{V}$ are of order $(\omega_0/m)^{1/2}$, so that the second-order term in $\mathbf{V}$ is of the same order of magnitude $(e^2\mathbf{a}^2/m)$ as the first-order term in W. That they actually cancel is seen most easily by making a gauge transformation of the type sometimes used to introduce the dipole approximation. This reduces $\mathbf{a}$ to zero while generating a new component

$$a'_\tau(\tau,\mathbf{x}) = -\mathbf{x}\cdot(d\mathbf{a}/d\tau) \quad .$$

In this gauge the coefficients c_n' satisfy the equation

$$i\,\frac{dc_n'(\tau)}{d\tau} = E_n\, c_n'(\tau) - e\,\frac{d\underset{\sim}{a}}{d\tau}\cdot\sum_{n'} \underset{\sim}{X}_{nn'}\, c_{n'}'(\tau)\,, \qquad (12)$$

with

$$\underset{\sim}{X}_{nn'} = \int d\zeta\, d^2\underset{\sim}{x}\;\; u_n^{+}(\zeta,\underset{\sim}{x})\,\underset{\sim}{x}\, u_{n'}(\zeta,\underset{\sim}{x})\,. \qquad (13)$$

This is an interesting formula because, although (12) looks like the familiar non-relativistic expression in the dipole approximation, it is in fact exact. [Recall that as compared with the usual integral, the integrand of (13) contains an extra factor of $e^{i(E_n-E_{n'})\zeta}$. This is just the standard exponential evaluated for an "energy-conserving" wave vector of magnitude $E_{n'} - E_n$ in the z direction.]

Since the matrix elements of X are of order $(m\omega_0)^{-1/2}$ and the diagonal matrix elements vanish, it is clear from (12) that the leading term in the energy-level shift is of order $(\omega/\omega_0)^2\,(e^2\underset{\sim}{a}^2/m)$, showing that the terms of order $(e^2\underset{\sim}{a}^2/m)$ must cancel exactly. As $\omega \to 0$ we obtain precisely the usual second-order Stark shift, as we must.

III. ELECTRON GREEN'S FUNCTION AND EFFECTIVE POTENTIALS

The electron Green's function in a plane wave has been calculated by numerous authors [4]. Here we want to put it in a particularly convenient form. In the two-component formalism, and using the notation (10), it may be written as

$$G(x,x') = \frac{i}{(2\pi)^3}\int d\eta\, d^2p\left[\theta(\eta)\,\theta(\tau-\tau') - \theta(-\eta)\,\theta(\tau'-\tau)\right]$$

$$e^{-i\eta(\zeta-\zeta')+i\underset{\sim}{p}\cdot(\underset{\sim}{x}-\underset{\sim}{x}')}\; \exp\left(-i\int_{\tau'}^{\tau} d\tau''\,\hat{E}_p(\tau'')\right). \qquad (14)$$

Note that in common with the wave functions, it has no explicit spin dependence.

One can extract the gauge dependence as a line integral factor [16)]

$$G(x,x') = e^{-ie\Lambda(x,x')}\,\hat{G}(x,x') \qquad (15)$$

with

$$\Lambda(x,x') = \int_{x'}^{x} dx^{\mu}\, a_{\mu}(x) = -(\underset{\sim}{x}-\underset{\sim}{x}')\cdot\langle\underset{\sim}{a}\rangle\;,$$

where the angular brackets denote a time average between τ' and τ:

$$\langle\underset{\sim}{a}\rangle = \frac{1}{\tau-\tau'}\int_{\tau'}^{\tau} d\tau''\,\underset{\sim}{a}(\tau'')\;.$$

The gauge-invariant factor $\hat{G}$ may be written:

$$\hat{G}(x,x') = \int \frac{d^4p}{(2\pi)^4}\;\frac{e^{-ip\cdot(x-x')}\,2\eta}{M^2(\tau,\tau') - p^2 - i\varepsilon}\;, \qquad (16)$$

where M^2 denotes the mass operator

$$M^2(\tau,\tau') = m^2 + e^2\langle\underset{\sim}{a}^2\rangle - e^2\langle\underset{\sim}{a}\rangle^2\;.$$

The non-Lorentz-invariant factor 2η in the integrand of (16) is required by the Lorentz transformation properties of $\psi(x)$. (Under longitudinal boosts, $\eta^{-1/2}$ is invariant.)

Because of the plane-wave character of the field, three components of energy momentum are conserved. In fact, the Fourier transform of (16) is

$$\hat{G}(p,p') = 2\eta(2\pi)^3\,\delta(\eta-\eta')\,\delta_2(\underset{\sim}{p}-\underset{\sim}{p}')$$

$$i\int d\tau\, e^{i(E-E')\tau}\int_0^{\infty} d\alpha\, e^{-i\alpha[M^2(\tau+2\alpha\eta,\tau) - p^2]}\;. \qquad (17)$$

A particularly simple and interesting case is that of a circularly polarized monochromatic wave

$$\underset{\sim}{a}(\tau) = \mathrm{Re}(\underset{\sim}{a} e^{-i\omega\tau}) \; , \qquad \underset{\sim}{a}^2 = 0 \; . \tag{18}$$

For this case a simple calculation shows that M^2 depends only on the difference of its arguments,

$$M^2(\tau,\tau') = m^2 + \Delta m^2 \left\{ 1 - \frac{\sin^2\left[\frac{1}{2}\,\omega(\tau-\tau')\right]}{\left[\frac{1}{2}\,\omega(\tau-\tau')\right]^2} \right\}$$

$$= M^2(\tau-\tau') \; , \tag{19}$$

say, where

$$\Delta m^2 = \frac{1}{2} e^2 \; \underset{\sim}{a}\cdot\underset{\sim}{a}^*$$

is the intensity-dependent mass shift. In this special case, $\hat{G}$ conserves all four components of energy momentum and can be written in a particularly compact form,

$$\hat{G}(p,p') = 2\eta \; (2\pi)^4 \; \delta_4(p-p') \; i \int_0^\infty d\alpha \; e^{-i\alpha[M^2(2\alpha\eta)-p^2]} \quad . \tag{20}$$

Physically, this means that the absorption and emission of photons from a circularly polarized plane wave is completely described by the time-integral factor in (15).

It is clear from (19) that while M^2 approaches the shifted mass $m^2 + \Delta m^2$ as $\tau - \tau' \to \infty$, for short time intervals it is approximately equal to the unshifted mass m^2 . In particular $G(x,x)$ is independent of a_μ .

Now let us consider the possible effect of a plane-wave field a_μ on the effective potential function $V(\varphi)$ for some neutral scalar field φ . There is only one way to couple zero-momentum φ lines to a charged-fermion line, which amounts to replacing the fermion mass m by $m + g\varphi$. In particular, the one-loop contribution to $V(\varphi)$ is obtained from the single vacuum loop by this replacement. Since the vacuum loop involves $G(x,x)$, the one-loop potential is independent of $\underset{\sim}{a}$.

As we noted in the introduction, this conclusion holds more generally, to all orders in the loop expansion, as a consequence of gauge and Lorentz invariance. However, it is true only for a plane-wave field, and only for vanishing external momenta. In other situations we may expect an effect.

Consider for example the one-loop contribution to the scalar field self-energy function π_φ (Fig.1), assuming the simplest invariant coupling of φ to the charged fermions. For the case of circular polarization (18), π_φ is translationally invariant and its Fourier transform is:

$$\pi_\varphi(q) = \frac{ig^2}{(2\pi)^4} \int d^4p \int d^4p' \; \delta(p + p' - q)$$

$$\int_0^\infty d\alpha \; e^{-i\alpha[M^2(2\alpha\eta)-p^2]} \int_0^\infty d\beta \; e^{-i\beta[M^2(2\beta\eta') - p'^2]} \; .$$

Using dimensional regularization to handle the quadratic divergence, the momentum integrals can be readily evaluated to give

$$\pi_\varphi(q) = -g^2 \int_0^\infty d\alpha \int_0^\infty d\beta \; \left[4\pi i(\alpha + \beta)\right]^{-n/2}$$

$$\exp\left[i\,\frac{\alpha\beta}{\alpha + \beta}\, q^2 - i(\alpha+\beta)\, M^2\!\left(2\,\frac{\alpha\beta}{\alpha + \beta}\, q_\zeta\right)\right] \; . \tag{21}$$

This formula exhibits π_φ as a function of the invariants q^2 , $q\cdot k = q_\zeta\omega$ and a^2 . Clearly, the coefficient of a^2 vanishes when q or k tend to zero like $(q\cdot k)^2$, as it should since the combination $(q\cdot k)^2\,a^2$ is the gauge-invariant quantity $(F_{\mu\nu}q^\nu)^2$. For a neutral particle at rest, the energy shift induced by the plane wave is therefore smaller than the electron mass shift by a factor ω^2/m^2 . Nevertheless, in suitable circumstances it might be observable.

Another case in which the external field could have an effect is that of a non-plane-wave field. The effective potential function $V(\varphi)$ must be a scalar function of $F_{\mu\nu}$, and can therefore have a non-trivial dependence in the external field when $F_{\mu\nu}F^{\mu\nu} \neq 0$. The contribution of a single plane wave vanishes essentially because there is only one wave vector k^μ involved, but when two or more such vectors appear this is no longer the case. In particular we may consider a superposition of plane waves travelling in opposite directions, described (in cartesian co-ordinates) by

$$k^\mu = (\omega,0,0,\omega) \quad \text{and} \quad k'^\mu = (\omega,0,0,-\omega).$$

Physically, such a wave is readily produced by reflecting a laser beam from a mirror. Unfortunately, diagrams such as Fig.2 with only two external photon lines, and any number of zero-momentum φ lines, will give no contribution because of the delta function restricting the photons to have equal and opposite energy-momentum vectors. However, diagrams with four external photon lines will give a contribution of order $e^2\omega^4 a^4$ to $V(\varphi)$. In present circumstances this contribution is far too small to have observable consequences, but this may not always be true.

IV. CONCLUSIONS

We have shown that the null result of the experimental search[6] for intensity-dependent level shifts in an atom in a plane-wave field was to be expected, despite the fact that the e^2A^2 interaction term is more properly regarded as generating a mass shift rather than a uniform energy shift. Indeed the shift is not uniform but is cancelled by second-order effects of the $ep\cdot A$ term.

Similarly we have shown that a pure plane-wave field will not induce restoration of spontaneously broken symmetries, because its contribution to the effective potential function vanishes for reasons of invariance. In principle, at least, a non-plane-wave field can have the desired effect. The simplest practical example is the standing wave produced by reflection from a mirror. However, at the present time the effect is far too small to be seen.

REFERENCES

1) N.D. Sengupta, Bull. Math. Soc. Calcutta 44, 175 (1952).

2) A.I. Nikishov and V.I. Ritus, Zh. Eksperim. i. Teor. Fiz. 46, 776 (1963), [English translation: Soviet Phys.-JETP 19, 529 (1964)].

3) L.S. Brown and T.W.B. Kibble, Phys. Rev. 133, A705 (1964).

4) J.H. Eberly, in Progress in Optics, Ed. E. Wolf (North-Holland 1969), Vol.7.

5) H.R. Reiss, Phys. Rev. Letters 17, 1162 (1966).

6) J.R. Mowat, C.E. Johnson, H.A. Shugart and V.J. Ehlers, Phys. Rev. D3, 43 (1971).

7) T.W.B. Kibble, Phys. Rev. 150, 1060 (1966).

8) E.P. Tryon, Phys. Rev. Letters 32, 1139 (1974).

9) Abdus Salam and J. Strathdee, ICTP, Trieste, preprint IC/74/140 (to be published in Nucl. Phys. B).

10) W. Becker, private communication.

11) R.A. Neville and F. Rohrlich, Nuovo Cimento 1A, 625 (1971).

12) J.D. Bjorken, J.B. Kogut and D.W. Soper, Phys. Rev. D3, 1382 (1971).

13) J.S. Bell and H. Ruegg, "Hydrogen atom on null plane and Melosh transformation", CERN preprint TH.1970 (1975).

14) D.M. Volkov, Z. Physik 94, 250 (1935).

15) I.I. Goldman, Phys. Letters 8, 103 (1964).

16) Ref.3, Appx.A.

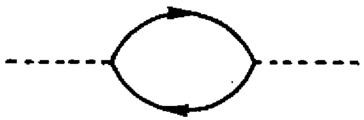

Fig.1 One-loop contribution to scalar-field self-energy.

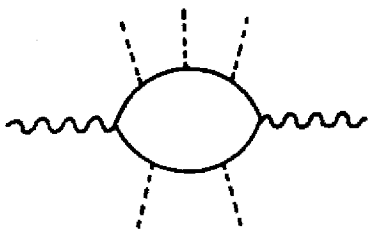

Fig.2 Photon self-energy modified by any number of zero-momentum φ lines.

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