as before. Once more, one cannot rule out the presence of this pole. Its existence would imply, in this case, that parity is not conserved in strong interactions.

The illustrations we have employed indicate that attempts to establish the presence or absence of certain vertex parts (and thus determine parities) by the above means are, at best, questionable. Of course, to make our point we have been selective in the choice of experiments. Nevertheless, this is in keeping with the spirit of the pole approximation as presently advocated.

The conjecture that poles in $\cos\theta$ appear in the S matrix is no doubt reasonable. However, an analysis of experiments which assumes that the poles predominate may require accurate measurements sufficiently close to the pole, yet far from other singularities. What is meant by "accurate," "sufficiently close," and "far" is somewhat nebulous. In fact, in the above illustrations [and in process (1)], in order to establish the presence or absence of a pole, it is necessary to distinguish between a possible zero near $\cos\theta = 1$ and a double zero at $\cos\theta = c_0$ (or α_0). At higher energies the poles approach the physical region, that is the points $\cos\theta = \pm 1$. Analysis of experiments at these higher energies probably would not eliminate the ambiguities, since the other singularities also approach these limit points.

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[1]S. Mandelstam, Phys. Rev. 112, 1344 (1958).

[2]J. G. Taylor, Nuclear Phys. 9, 357 (1959); J. G. Taylor (to be published).

[3]G. F. Chew, Phys. Rev. 112, 1380 (1958); G. F. Chew and F. E. Low, University of California Radiation Laboratory Report UCRL-8427 (unpublished); Taylor, Moravcsik, and Uretsky, Phys. Rev. 107, 563 (1957); M. Moravcsik, Phys. Rev. Letters 2, 352 (1959); S. Barshay and S. Glashow, Phys. Rev. Letters 2, 371 (1959).

[4]Ashkin, Glaser, Feiner, and Stern, Phys. Rev. 105, 724 (1959).

[5]The corresponding circle of convergence for process (1) lies entirely outside of the physical region. Nevertheless, Taylor[2] has suggested using data in the region $1 \geq \cos\theta \geq 0$.

ERRATA

INTEGRATION OF SECONDARY CONSTRAINTS IN QUANTIZED GENERAL RELATIVITY. P. W. Higgs [Phys. Rev. Letters 1, 373 (1958)].

In this Letter[1] it was stated that three of the secondary constraints[2] in the quantized theory of gravitation are equivalent to the statement that the wave functional $\Psi\{g_{rs}(x)\}$ is invariant under general transformations of the three coordinates x. From this the conclusion was drawn that Ψ must depend only on the three eigenvalues of the Ricci tensor $R_{rs}(x)$. (As before, Latin indices run from 1 to 3; Greek indices from 0 to 3.) It appears on more careful examination of the transformation properties of Ψ that the former statement is not quite correct, so the conclusion is erroneous. The purpose of the present note is to clarify the invariance properties of Ψ and to present the correct solution of the secondary constraints

$$\mathcal{H}_u \Psi = 0, \tag{1}$$

where

$$\mathcal{H}_u = g_{rs,u}\pi^{rs} - 2(g_{ru}\pi^{rs})_{,s} \quad \text{with} \quad \pi^{rs} = -i\delta/\delta g_{rs}$$

Let us consider just the localized transformations of the three spatial coordinates, $x'^r = x^r + a^r(x)$, where the infinitesimal functions $a^r(x)$ vanish at infinity. Then

$$\delta g_{rs}(x) = -(a^u g_{rs,u} + a^u{}_{,r} g_{us} + a^u{}_{,s} g_{ru}),$$

and so, the primary constraints having been satisfied already,

$$\delta\Psi = \int d^3x\, \delta g_{rs}(x)\delta\Psi/\delta g_{rs}(x)$$
$$= -i\int d^3x\, a^u(x)\mathcal{H}_u(x)\Psi - 2i\int dS_r\, a^u(x) g_{us}\pi^{rs}\Psi.$$

The first term vanishes if the constraint (1) is satisfied; the second is an integral over the surface at infinity which vanishes only if $a^u(x)$ is localized. We therefore have to construct functionals $\Psi\{g_{rs}(x)\}$ which are invariant only under localized transformations.

The problem is solved by transforming g_{rs} into a form which satisfies certain transversality conditions:

$$g_{rs}(x) = y^a{}_{,r}y^b{}_{,s}g^T{}_{ab}(y).$$

If $y^a(x)$ are three independent functions such that the differences $y^a - x^a$ are localized and $g^T{}_{ab}$ satisfies suitable conditions, then we know that Ψ is independent of y^a. By suitable conditions we mean that $g^T{}_{ab}$ has only three independent components and the functions y^a are uniquely determined by $g_{rs}(x)$ together with these conditions. It turns out that the well-known harmonic coordinate conditions fulfill these criteria.

It is convenient to introduce the contravariant densities h^{rs}, related to g_{rs} by $h^{rs}g_{st} = (-g)^{1/2}\delta^r{}_t$. Then we define the "transverse gravitational potentials"

$$h^{Tab}(y) = (\det y^c{}_{,t})^{-1} y^a{}_{,r} y^b{}_{,s} h^{rs}(x), \qquad (2)$$

where $y^a(x)$ is the regular solution of the generalized Laplace equation

$$(h^{rs}y_{,r})_{,s} = 0 \qquad (3)$$

which satisfies the asymptotic condition (localization)

$$\lim_{|x|\to\infty}(y^a - x^a) = 0\ .$$

The functions y^a and consequently also the transverse potentials are uniquely determined by $g_{rs}(x)$. The constraints (1) now tell us that $\Psi = \Psi\{h^{Tab}(y)\}$. The transversality condition, which follows from (2) and (3), is

$$h^{Tab}{}_{,b} = 0, \qquad (4)$$

so there are only three independent transverse components. They may be exhibited explicitly by performing a Fourier transformation and resolving the transform of h^{Tab} parallel and perpendicular to the wave vector k_a: we define

$$\gamma^{ij}(k) = \int d^3y\, \exp(-ik\cdot y) e^i{}_a e^j{}_b[h^{Tab}(y) + \delta^{ab}], \qquad (5)$$

where $e^i{}_a(k)$ is an orthonormal triad of vectors (the scalar product here is Euclidean) such that $k_a = |k| e^3{}_a$. Then, by (4) and (5), the only nonvanishing components of γ^{ij} are γ^{11}, γ^{22}, $\gamma^{12} = \gamma^{21}$. [In addition, the reality condition $\gamma^{ij}(-k) = \gamma^{*ij}(k)$ is to be imposed.] Finally, the wave functional may be written as $\Psi\{\gamma^{ij}(k)\}$.

The transverse potentials which have been defined here reduce in the linearized theory to those which were obtained by Arnowitt and Deser,[3] except that their γ^{ij} was traceless on account of the eighth constraint. It remains to be seen whether Dirac's $\mathcal{H}_L$ condition implies a similar restriction in the full theory.

[1]Equation (11) of this paper should read

$$[\mathcal{H}_u(x), \mathcal{H}_L(y)] = i\mathcal{H}_L(x)\partial/\partial x^u \delta^{(3)}(x-y).$$

[2]P. A. M. Dirac, Proc. Roy. Soc. (London) A246, 333 (1958).

[3]R. Arnowitt and S. Deser, Phys. Rev. 113, 745 (1959).

EFFECT OF NUCLEAR ELECTRIC DIPOLE MOMENTS ON NUCLEAR SPIN RELAXATION IN GASES. P. A. Franken and H. S. Boyne [Phys. Rev. Letters 2, 422 (1959)].

In this Letter it was estimated that careful measurements of nuclear spin relaxation times in noble gases such as He^3 and Xe^{129} could reveal the existence of nuclear electric dipole moments as small as 10^{-3} or 10^{-4} nuclear magneton. (One nuclear magneton $= eh/2Mc = e\times 10^{-14}$ cm.) Indeed, we interpreted the already available measurements on these gases to indicate upper limits of order 10^{-1} to 10^{-2} nuclear magneton.

Professor E. M. Purcell has brought to our attention that we have made a serious error in