of Eq. (15) becomes a differential equation for ν_1 :

$$(1-u^2)\frac{d^2\nu_1}{du^2}-5u\frac{d\nu_1}{du}+3(\lambda-1)\nu_1=0. \quad (21)$$

To ensure that Eq. (5) is satisfied, we must impose upon Eq. (21) the boundary condition that ν_{2j+1} , the $2j$ th derivative of ν_1 [see Eq. (20)], is zero. The required solution is then a polynomial in u :

$$\begin{aligned} \nu_1 &= \sum_r b_{2r-1} u^{2r-1}, \\ 2r(2r+1)b_{2r+1} &= [(2r-1)(2r+3) \\ &\quad - (2j-1)(2j+3)]b_{2r-1}, \end{aligned} \quad (22)$$

where the sum over r runs over $\frac{1}{2}, \frac{3}{2}, \dots, j$ when j is a half-integer, and over $1, 2, \dots, j$, when j is an integer. We can now use Eq. (22) in conjunction with Eqs. (18) and (19) to determine those functions f and g that yield a linear realization (j, j) constructed from the pion field alone. For example, in the special cases $j=\frac{1}{2}$ and 1 we find that

$$\begin{aligned} j=\frac{1}{2}: \quad f^2+\pi^2 &= \alpha, & g &= 0, \\ j=1: \quad f^2-2\alpha f+\pi^2 &= 0, & g &= -1/f, \end{aligned} \quad (23)$$

where α is an arbitrary constant.

It is easy to show from the commutation rules in Eq. (2) that $h_n(\pi_+)^n$ [see Eq. (16)] is an eigenstate of isospin with $T=T_3=n$ ($n=0, 1, \dots, 2j$). Thus, by operating on these states with the isospin-lowering operator $T_- = T_1 - iT_2$, we can generate the complete set of isospin states contained in the representation (j, j) , and hence the complete representation itself.

Now suppose that we have a new isovector field $\tilde{\pi}_a$ for which

$$K_a \tilde{\pi}_b = -i(\delta_{ab}F + \tilde{\pi}_a \tilde{\pi}_b G), \quad (24)$$

where F and G are functions of $\tilde{\pi}^2$. In order to convert this to a linear realization, we need to transform $\tilde{\pi}_a$ into π_a and F and G into the functions f and g of the preceding discussion. Following Weinberg,⁵ we seek a redefinition of the form

$$\begin{aligned} \pi_a &= \tilde{\pi}_a \Phi(\tilde{\pi}^2), & f &= F\Phi, \\ g &= [G\Phi + 2(F + \tilde{\pi}^2 G)\Phi']/\Phi^2. \end{aligned} \quad (25)$$

If such a redefinition is to exist for any F and G , then we must be able to express Φ in terms of $\tilde{\pi}^2$ and F .

This expression is not difficult to find. If we define quantities $\tilde{\sigma}$ and $\tilde{u}$ analogous to σ and u in Eq. (19), then we find that

$$\sigma = \tilde{\sigma}\Phi, \quad u = \tilde{u}. \quad (26)$$

The condition $h_1(\pi^2) = -\lambda$ of Eq. (18) then leads to

$$\Phi = -(1/\lambda\tilde{\sigma})\nu_1(\tilde{u}). \quad (27)$$

Thus Φ in Eq. (27) is the redefinition required to transform the nonlinear realization of Eq. (24) into a linear one of the type (j, j) . We also note that

$$(\pi_+)^n h_n(\pi^2) \equiv (\tilde{\pi}_+)^n h_n(\tilde{\pi}^2), \quad (28)$$

and hence we obtain all the states of this representation. Since j has been left arbitrary in this discussion, it follows that our procedure can be used to convert a given nonlinear realization into any one of the allowed linear realizations.

As a final point, we observe that the quantities in Eq. (28) satisfy the integrability condition⁶ required by Weinberg for the existence of the matrix $\Lambda(\pi)$. It follows that the components of the matrix are determined by these quantities.

⁵ S. Weinberg, Ref. 1, Eqs. (2.12)–(2.15).

⁶ S. Weinberg, Ref. 1, Eq. (5.6).

Errata

Inelastic π^+n Interactions in the Center-of-Mass Energy Range 1.40–1.65 GeV, P. J. LITCHFIELD [Phys. Rev. **183**, 1152 (1969)]. The Hulthén form for the distribution of the spectator proton momentum was misquoted. It should read

$$P_p \propto \left(\frac{1}{\alpha^2 + p^2} - \frac{1}{\beta^2 + p^2} \right)^2 p^2.$$

The correct form was used in the analysis.

The Gravitational Field of a Disk, THOMAS MORGAN AND LESLEY MORGAN [Phys. Rev. **183**, 1097 (1969)]. Equation (26) should read

$$\phi = \frac{1}{2} \ln \{ (\omega/2b) [1 + (1 + 4b^2 \rho^2)^{1/2}] \}.$$