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

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

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:

$$\nu_1=\sum_r b_{2r-1}u^{2r-1},$$
$$2r(2r+1)b_{2r+1}=[(2r-1)(2r+3)-(2j-1)(2j+3)]b_{2r-1}, \tag{22}$$

where the sum over r runs over $\frac{1}{2}, \frac{3}{2}, \ldots, j$ when j is a half-integer, and over 1, 2, . . . 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

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

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 $T=T_3=n$ $(n=0, 1, \cdots, 2j)$. 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

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

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

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

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

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

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

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

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

$$(\pi_+)^nh_n(\pi^2)\equiv(\tilde{\pi}_+)^nh_n(\tilde{\pi}^2), \tag{28}$$

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 condition[6] 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

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

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

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