Refs. 5 and 6 are in disagreement with those of Ref. 3 (as well as with one another). If one of these recent calculations is correct, then the mass correction increases more rapidly with the magnetic field strength than Demeur's, and at fields of the order of 10^{13} G or larger it can no longer be assumed to be small.

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Comments on Isospin-Factored Current Algebra at Infinite Momentum

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A solution of the isospin-factored current algebra at infinite momentum was given by Chang, Dashen, and O'Raifeartaigh, who assumed that one could project from the isospin-factored current algebra those solutions which contain a spacelike part which is definitely coupled to a timelike part by the current. This assumption was used to conclude that the operator M_3 vanishes. However, since the resulting solutions still contain a spacelike part which is coupled to the timelike part, it was suspected that this assumption may have excluded an important class of solutions. This, however, is shown not to be the case.

I. INTRODUCTION

Recently a solution of the $SU(2)\otimes SU(2)$ isospin-factored current algebra at infinite momentum was presented by Chang, Dashen, and O'Raifeartaigh[1] in which it was discovered by using an angular condition that the resulting solutions were either physically trivial or that they contained a spacelike part, i.e., $M^2<0$, for the mass operator. The physically trivial solutions are those which correspond to the charge-density current-density algebra and to the free quark model which has the special property that the spacelike and timelike solutions are unconnected by the current. The existence of a spacelike part in other cases of solutions is undesirable, for it makes it very difficult to see how one could carry out a program of saturating the current-algebra equations at infinite momentum.

The solutions to the isospin-factored current algebra were obtained subject to the assumption that the current involved does not connect the spacelike and timelike states, expressed technically as the vanishing of the operator M_3 which is defined later. Since, in fact, it is found that the solutions do indeed contain a spacelike part, there existed the possibility that an important class of solutions had been excluded as a result of this assumption. It is shown, however, that the assumption regarding the connecting of spacelike and timelike states is actually unnecessary so that no solutions have been excluded.

In order to find solutions to the current algebra at infinite momentum, it is necessary to introduce an angular condition[2,3] so that one can generate the Lorentz group, for as a result of transforming to an infinite-momentum frame, the Lorentz group degenerates to an $E(2)\otimes D$ group,[4] i.e., the semidirect product of the group of Euclidean motions in a plane and a boost in a fixed direction. For the isospin-factored algebra, this angular condition can be reduced to a simple set of equations from which one can find various expressions for the mass operator. Central to the problem of finding solutions to this set of equations is the demonstration of the existence of an $E(2)\otimes D$ subgroup of the $SL(2C)$ groups which are generated by the basic operators. It is also shown in this work that the existence of the Lie algebra corresponding to this $E(2)\otimes D$ group is sufficient to make unnecessary the assumption that the current does not connect spacelike and timelike states.

II. ANGULAR CONDITION

The isospin algebra at infinite momentum has the form[5]

$$[F^a(\mathbf{k}),F^b(\mathbf{k}')]=i\epsilon_{abc}F^c(\mathbf{k}+\mathbf{k}'), \tag{2.1}$$

where $\mathbf{k}$ is the transverse momentum and where ϵ_{abc} is the Levi-Civita symbol. If one makes the assumption that the current is factored into the product of an isotopic spin generator of the Lie algebra of $SU(2)$ and a reduced matrix element such that

$$F^a(\mathbf{k})=I^aF(\mathbf{k}), \tag{2.2}$$

then the reduced matrix element must satisfy the equation

$$F(\mathbf{k})F(\mathbf{k}')=F(\mathbf{k}+\mathbf{k}'). \tag{2.3}$$

This equation forms a unitary representation of the

[1] S. J. Chang, R. Dashen, and L. O'Raifeartaigh, Phys. Rev. **182**, 1819 (1969).
[2] S. J. Chang, R. Dashen, and L. O'Raifeartaigh, Phys. Rev. **182**, 1805 (1969).
[3] R. F. Dashen and M. Gell-Mann, Phys. Rev. Letters **17**, 340 (1966).
[4] S. J. Chang and L. O'Raifeartaigh, J. Math. Phys. **10**, 21 (1969).
[5] M. Gell-Mann, in *Proceedings of the Erice Summer School, 1966* (Academic, New York, 1966).

Abelian group, and by a theorem due to Stone[1] has the general solution

$$F(\mathbf{k}) = e^{i\mathbf{k}\cdot\mathbf{x}}, \tag{2.4}$$

where $\mathbf{x} = (x_1, x_2)$ are two self-adjoint commuting operators. Not all $\mathbf{x}$ are allowed as solutions to the isospin-factored current algebra since Lorentz covariance requires only those which satisfy an angular condition. That is, one must find the necessary and sufficient condition for an operator $J(0)$, defined by

$$J(0) = \lim_{k\to\infty} e^{-i\sinh^{-1}kK_3} J_0(0) e^{i\sinh^{-1}kK_3} = J_0(0) + J_3(0), \tag{2.5}$$

where K_3 is the generator for a boost in the z direction, so that the operator $J_\mu(0)$ transforms as a current under the Lorentz group L. This problem has been solved, and it is found that $J(0)$ must transform as a scalar under the subgroup $E(2)$ of $SL(2C)$. This subgroup is generated by $E_1 = K_1 + L_2$, $E_2 = K_1 - L_2$, and L_3, where $\mathbf{L}$ and $\mathbf{K}$ are, respectively, the usual generators of $SL(2C)$ corresponding to rotations and accelerations. If one now introduces an additional rotation about any axis other than the z axis, it is possible with the $E(2)$ subgroup to generate the full Lorentz group L. This is the so-called angular condition, and it has been shown to be equivalent to the relations

$$[J_3, F(\mathbf{k})] = i\epsilon_{3ij} k_i \frac{\delta}{\delta k_j} F(\mathbf{k}), \tag{2.6}$$

$$\{I,\{I,\{I,F(\mathbf{k})\}\}\} = \{I,\{J,F(\mathbf{k})\}\}, \tag{2.7}$$

with

$$2\{I,\Theta\} = \delta(M^2)\delta(J_3)\Theta + [\mathbf{k}\cdot 2M\mathbf{J},\Theta] - k^2[J_3,\Theta]_+, \tag{2.8}$$

$$4\{I,\Theta\} = \delta^2(M^2)\Theta + 2k[M^2,\Theta]_+ + k^4\Theta, \tag{2.9}$$

where $\mathbf{J}$ and M are, respectively, generators for helicity transformations and the mass operator acting on states in the space $\mathcal{H}_\infty$ which results from the infinite-momentum limit. In the above expressions, use has been made of the notation

$$\delta^0(A)\Theta = \Theta, \quad \delta^1(A)\Theta = [A,\Theta], \quad \delta^2(A)\Theta = [A,[A,\Theta]], \text{ etc.}$$

III. REDUCTION OF ANGULAR CONDITION

The procedure which is followed to reduce the angular condition to a set of k-independent equations is to multiply (2.7) by $e^{-i\mathbf{k}\cdot\mathbf{x}}$ and expand in powers of k. A simplified set of equations can be obtained if one makes the assumption that the function can be continued analytically to the value $\mathbf{k} = (k, \pm ik)$, so that $k^2 = 0$. This assumption is physically reasonable since the function $F(\mathbf{k})$ corresponds to an electromagnetic form factor. It has been shown that in this case the angular condition is equivalent to the k-independent equation

$$M_{\pm\pm\pm} = \delta^3(X_\pm)M^2 = 0, \tag{3.1}$$

$$[X_\pm, \Lambda_\pm] = (\Lambda_\pm)_\pm = 0, \tag{3.2}$$

where

$$X_\pm = X_1 \pm iX_2 \tag{3.3}$$

and

$$\Lambda_\pm = \pm 2iMJ_\pm + \tfrac{1}{2}[M^2, X_\pm]_+. \tag{3.4}$$

Use is now made of the information found for the case when $k^2 = 0$ to study the situation for other values of k^2. It is convenient to introduce the vector $\mathbf{k} = (k,0)$ when expanding (2.7) in powers of k. But first it is useful to write the condition that the helicity operator $\mathbf{J}$ commutes with the mass operator M^2 and the condition that J_+, J_-, and J_3 generate the algebra $SU(2)$, in the form

$$[\Lambda_\pm, M^2] = \tfrac{1}{2}[M^2, M_\pm]_+, \tag{3.5}$$

$$[\Lambda_+, \Lambda_-] = 4[M^2, B]_+ - \tfrac{1}{4}[M_+, M_-] \tag{3.6}$$

with

$$B = 2J_3 + \tfrac{1}{8}[(\Lambda_-)_+ - (\Lambda_-)_+], \tag{3.7}$$

where

$$(\Lambda_\mp)_\pm = \delta(X_\pm)\Lambda_\mp.$$

For an arbitrary operator Θ, (2.8) can be written as

$$2e^{-i\mathbf{k}\cdot\mathbf{x}}\{I, e^{i\mathbf{k}\cdot\mathbf{x}}\Theta\} = \sum_{n=0}^{\infty} (-i)^n G_n(\Theta) k^n, \tag{3.8}$$

with

$$G_n(\Theta) = M_n\delta(J_3)\Theta - inM_{n-1}X_2\Theta - (J_3)_n\Theta M^2 + 1_n\Theta M^2 J_3 - in[(2MJ_1)_{n-1}\Theta - 1_{n-1}\Theta 2MJ_1] + n(n-1)[(J_3)_{n-3}\Theta + 1_{n-2}\Theta J_3], \tag{3.9}$$

where use is made of the notation $\delta^n(X_1)\Theta \equiv \Theta_n$ for the commutator of Θ taken n times with X_1, but we have used the notation M_n to mean $\delta^n(X_1)M^2$. In a similar manner (2.9) can be expanded in powers of k, and it becomes

$$4e^{-ikx_1}\{J, e^{ikx_1}\} = \sum_{n=0}^{\infty} \frac{(-i)^n}{n!} E_n k^n, \tag{3.10}$$

with

$$E_n = (M^4)_n - 2M_n M^2 + 1_n - n(n-1)2(M_{n-2} + 1_{n-2}M^2) + n(n-1)(n-2)(n-3)1_{n-4}. \tag{3.11}$$

For the special case when $\Theta = 1$, one finds for (3.8)

$$2e^{-ikx_1}\{I, e^{ikx_1}\} = \sum_{n=0}^{\infty} (-i)^n \frac{a_n}{n!} k^n, \tag{3.12}$$

with

$$a_n = G_n(1), \quad a_0 = 0, \quad a_1 = i\delta(X_2)M^2, \quad a_n = 2na_{n-2} \quad \text{for } n \geq 2, \tag{3.13}$$

with B given by (3.7) and with C given by

$$C = \tfrac{1}{4}\delta(X_1)\delta(X_2)M^2. \tag{3.14}$$

It should also be observed that B and C are self-adjoint operators on the space $\mathcal{H}_\infty$ and that $C_2 \equiv \delta^2(X_1)C=0$, and $A_n = B_n$ for $n \geq 2$ as a result of (3.1).

In continuing the reduction process for the angular condition, we now introduce the difference expression

$$4e^{-ikx_1}(\{I,\{I,e^{ikx_1}\}\}-\{J,e^{ikx_1}\}) = \sum_{n=0}^{\infty} \frac{(-i)^n}{n!} D_n k^n, \quad (3.15)$$

with

$$D_n = \sum_{l=0}^{\infty} G_2(a_{n-l}) - \frac{E_n}{n!}, \quad (3.16)$$

and the difference expression

$$8e^{-ikx_1}(\{I,\{I,\{I,e^{ikx_1}\}\}\}-\{I,\{J,e^{ikx_1}\}\}) = \sum_{p=0}^{\infty}\sum_{q=0}^{\infty} G_q(D_{p-q})(-i)^p \binom{p}{q} \frac{1}{p!} k^q, \quad (3.17)$$

which gives for each order p of the momentum k the desired k-independent set of equations for the angular condition (2.7), namely, the set of equations

$$\sum_{q=0}^{p} G_q(D_{p-q}) \binom{p}{q} = 0. \quad (3.18)$$

This set of equations would not be of much use if it did not terminate for some finite value of p. The termination of the series (3.17) depends upon the vanishing of the basic operators M^2, $\mathbf{J}$, B, and C as a result of applying the commutative operation $\delta^m(X_+)\delta^n(X_-)$ to them, where m and n are integers. It is now straightforward to show that

$$B_{+-}=0, \quad (3.19)$$

where

$$B_\pm = \delta(X_\pm)B = X_\pm \pm \tfrac{1}{8}(\Lambda_\mp)_{\pm\pm}. \quad (3.20)$$

In continuing the study of the termination question, we will now show that $M_4=0$. To prove this, we observe firstly that M_n can be written as

$$M_n = \frac{1}{2^n} \sum_{m=0}^{n} \binom{n}{m} \delta^{n-m}(X_+)\delta^m(X_-)M^2 \quad (3.21)$$

so that

$$M_4 = \frac{1}{2^4}\binom{4}{2} M_{++--}. \quad (3.22)$$

Upon applying the commutation operator $\delta^3(X_+)$ to both sides of (3.5), one finds the important result

$$(M_{++})^2=0. \quad (3.23)$$

If one now applies $\delta^4(X_-)$ to the equation using the property that $M_5=0$, it is found that

$$(M_4)^2=0. \quad (3.24)$$

Since M_4 is Hermitian, one has the result that it vanishes. Finally, we conclude that the termination of (3.17) is dependent upon the existence of some integer n such that $\delta^n(X_1)B=0$.

It has not been proved in general that the commutator of B taken n times with X_1 does in fact vanish; however, if B is expressed as a nonlinear combination of the basic operators $X_\pm$, $M_\pm$, $M_{\pm\pm}$, and $M_{\pm\pm\pm}$ such that $\delta(J_3)B=0$, then it is seen that B_3 vanishes. On the other hand, if B is represented in terms of Λ_+ and Λ_-, then one finds a commutation equation such that $\delta^n(X_+)B$ depends upon $\delta^{n-1}(X_+)B$. The Hermitian operator B must also satisfy the following:

$$[B_{+++},M_{--+}]+4[B_{++},M_{-+}]=0, \quad [B_{--},B_{++}]=0. \quad (3.25)$$

The first of these expressions is obtained by applying $\delta^4(X_+)\delta^4(X_-)$ to (3.5), and the second is found by applying $\delta^3(X_-)\delta^3(X_+)$ to (3.6).

At this point it is assumed that the expansion (3.17) does in fact terminate. This assumption is made reasonable by the fact that the solutions for B that are found are independent of the operators Λ_+ and Λ_- so that they do vanish as a result of forming the commutator of B a finite number of times with X_+. A consequence of this assumption is that

$$A_2=B_2=0. \quad (3.26)$$

This can be seen by observing that if (3.17) terminates for some power $p=(N)^3$, then

$$(a_N)^3=(B_{N-2})^3=0 \quad (3.27)$$

for all $N\geq 4$. Since B_{N-2} is either Hermitian or anti-Hermitian, we obtain (3.26). Furthermore, from (3.19) and $\delta(J_3)B_{\pm\pm}$ we obtain the result

$$B_{\pm\pm}=0. \quad (3.28)$$

The next step in the reduction is to establish the important result

$$M_3=0. \quad (3.29)$$

This result was used in Ref. 1 as a consequence of the assumption that one rejects solutions for which a spacelike part exists and is definitely coupled to a timelike part by the current e^{ikX_1}. To understand this, it is observed that if one performs an expansion in powers of k of the expectation value of the operator M^2 between the states $e^{ikX_1}|f\rangle$, then it is seen that the operator M_3 is associated with the leading term k^3, so that for large values of momentum k the expectation value of M^2 on the above states could be negative if M_3 is not set equal to zero. Since the final solutions of the isospin-factored current algebra do in fact contain spacelike parts it remained possible that an important class of solutions was lost by this assumption. However, this is now shown not to be the case since it is shown that (3.29) is a consequence of previously derived results. To establish this result, one begins by applying the operator $\delta^3(X_+)\delta^2(X_-)$ to (3.5) to find

$$[(\Lambda_-)_{+++},M_{---}]+3[(\Lambda_-)_{++},M_{--+}] =\tfrac{9}{2}[M_{--+},M_{-++}]_+. \quad (3.30)$$

Next use is made of (3.7) to show that

$$(\Lambda_-)_{+++}=0, \quad (3.31)$$

so that one has

$$[B_+,M_{+--}]=\tfrac{3}{16}[M_{+--},M_{++-}]_+. \quad (3.32)$$

If one now multiplies this equation on both sides from the left and also from the right with M_{+--}, one will find the two equations

$$M_{+--}B_+M_{+--}=\tfrac{3}{16}M_{+--}M_{-++}M_{+--}, \quad (3.33)$$

$$M_{+--}B_+M_{+--}=-\tfrac{3}{16}M_{+--}M_{-++}M_{+--}, \quad (3.34)$$

which are satisfied only if

$$(M_3)^3=M_{--+}M_{++-}M_{--+}=0. \quad (3.35)$$

Since M_3 is an anti-Hermitian operator, one finds the result given in (3.29). In obtaining this result, use has been made of the relation

$$(M_{+--})^2=0, \quad (3.36)$$

which is a consequence of applying $\delta^2(X_+)$ to the adjoint of (3.23). Furthermore, since

$$\delta(J_3)M_{\pm\pm}=\pm M_{\pm\pm}, \quad (3.37)$$

one finds

$$M_{\pm\pm}=0. \quad (3.38)$$

An important consequence of the vanishing of M_3 is the vanishing of B_1. To see this, one first observes that (3.19) and (3.38) require

$$C_1=0. \quad (3.39)$$

Also it is easily seen that (3.39) gives

$$G_n(\Theta)=0, \quad n\geq 3. \quad (3.40)$$

It is now seen that the highest remaining power in the expansion (3.17) is at order k^9, where it is found that

$$(a_3)^3=(B_1)^3=0. \quad (3.41)$$

But again it is observed that B_1 is anti-Hermitian and satisfies the eigenvalue equation

$$\delta(J_3)B_\pm=\pm B_\pm, \quad (3.42)$$

so that one concludes that

$$B_1=B_+=B_-=0. \quad (3.43)$$

The further reduction of the angular condition can now be accomplished by separating the quantities that occur in (3.18) according to the eigenvalues of the equation $\delta(J_3)\Theta_{(n)}=n\Theta_{(n)}$. After considerable algebra, one finds the remaining nonvanishing quantities occurring in (3.16) to be

$$D_{2(0)}=4[iK,M^2]+8M^2, \quad (3.44a)$$

$$\frac{1}{3!}D_{3(\pm1)}=[iK,\mathfrak{F}_+]+\mathfrak{F}_\pm, \quad (3.44b)$$

$$D_{4(0)}=4![4B^2-(R+1)^2], \quad (3.44c)$$

where K, $\mathfrak{F}_\pm$, and R are given by

$$iK\pm(J_3-B)=-\tfrac{1}{4}(\Lambda_\mp)_\pm, \quad (3.45)$$

$$\mathfrak{F}_\pm=\Lambda_\pm+M_{1(\pm1)}, \quad (3.46)$$

$$R=-\tfrac{1}{2}M_{2(0)}. \quad (3.47)$$

As an example of the technique used in deriving the above equations, one obtains after applying $\delta(X_-)$ to (3.5) the equation

$$[\Lambda_+,M_-]+[(\Lambda_+)_-,M^2] =\tfrac{1}{2}[M_{+-},M^2]+2[M_{1(1)},M_{1(-1)}]_+, \quad (3.48a)$$

which can be used to give (3.44a).

The final step in obtaining the k-independent equations corresponding to the angular condition is to use (3.44) in (3.17) and separate the equation according to the eigenvalues of $\delta(J_3)$. The algebraic procedure is similar to that which has already been used and will not be presented in detail. The finite set of k-independent equations corresponding to (3.17) can be written as

$$\text{order } k^6, \quad BG=0, \quad (3.48b)$$

$$\text{order } k^5, \quad M_{2(\pm2)}G_\pm=0, \quad (3.48c)$$

$$[\mathfrak{F}_\pm,G]-2M_{1(\pm1)}G+[\mathfrak{F}_\pm,G_\pm] -[X_\pm,G_\pm]M^2=0, \quad (3.48d)$$

$$\text{order } k^4, \quad M_{2(\pm)}G_0-4M_{1(\pm1)}G_\pm+[\mathfrak{F}_\pm,G_\pm] -[X_\pm,G_\pm]M^2=0, \quad (3.48e)$$

where

$$G_0=\tfrac{1}{8}D_{2(0)}, \quad (3.49a)$$

$$G_\pm=-\tfrac{1}{4}D_{3(\pm1)}, \quad (3.49b)$$

$$G=\frac{1}{4^2 3!}D_{4(0)}. \quad (3.49c)$$

IV. IMPLICATIONS OF $E(2)\otimes D$ SUBGROUP

A very important result which is used in finding mass operators which satisfy (3.48) is

$$[iK,\mathbf{X}]=\mathbf{X}. \quad (4.1)$$

This equation together with the equations

$$[X_+,X_-]=0, \quad (4.2a)$$

$$[J_3,K]=0, \quad (4.2b)$$

$$[J_3,X_\pm]=X_\pm \quad (4.2c)$$

generates an algebra isomorphic to the algebra of $E(2)\otimes D$, where $E(2)$ is the group of Euclidean motions in a plane, and D is a one-parameter boost in the z direction.

It is also possible to show that if one requires initially that this algebra is satisfied, then a necessary condition is that M_3 vanishes, provided one uses (3.1) and (3.2) which result from the case when k^2 is zero. First, we observe that (4.1) used with (3.2) and K written in the form

$$K=\tfrac{1}{8}[(\Lambda_-)_+ + (\Lambda_+)_-] \tag{4.3}$$

implies

$$X_+=\tfrac{1}{8}(\Lambda_-)_{++}. \tag{4.4}$$

This result together with the equation found by applying $\delta(X_+)$ to (3.4) gives

$$B_1=0. \tag{4.5}$$

Next we show that this result requires M_3 to vanish. To do this, we apply the commutation operation $\delta^2(X_+)\delta^2(X_-)$ to (3.6) and use (4.4) and its Hermitian adjoint to show that

$$[M_{3(1)},M_{3(-1)}]=0. \tag{4.6}$$

Upon multiplying this equation by $M_{3(1)}$ and using the property that $M_{3(1)}{}^2$ vanishes, one concludes that

$$(M_3)^3=0. \tag{4.7}$$

But since M_3 is an anti-Hermitian operator, one obtains the desired result.

In addition, if the condition (3.48a) is strengthened to the three separate conditions

$$B=0 \tag{4.8a}$$

$$=\pm\tfrac{1}{2}(R+1), \tag{4.8b}$$

and if one accepts the results of the case when k^2 vanishes, then it is easily seen that the conditions (4.8) require M_3 to vanish and therefore the preservation of the $E(2)\otimes D$ structure. The case $B=0$ is trivial since in this case B_+ is obviously zero, and we have already seen that this is sufficient to show that M_3 vanishes. For the case (4.8b), one finds

$$B_\pm=\mp M_{3(\pm1)} \tag{4.9}$$

so that (3.1) gives

$$B_2=0. \tag{4.10}$$

However, we have already seen in Sec. III that this is sufficient to prove that M_3 vanishes.

V. REVIEW OF SOLUTIONS

The mass spectra representing the solutions of the isospin-factored current algebra have been presented elsewhere,[1] and the details for the construction of the solutions can be found there; however, for completeness, a review of the results is presented here. In order to obtain solutions which are consistent with the k-independent equation (3.48), one finds solutions of the Hilbert spaces $\mathcal{H}_0$ and $\mathcal{H}_\pm$ which form a basis for the mass operator corresponding to the cases in (4.8). It can be shown that a general solution can be obtained by a suitable coupling of the solution for the three separate Hilbert spaces.

For each of the values of B, it is possible with a suitable redefinition of the basic operators $X_\pm$, K, $\mathbf{J}$, and $\mathcal{F}_\pm$ to generate an algebra isomorphic to the algebra of $SL(2C)$. These algebras can be used to construct mass operators which are consistent with the angular condition (3.48). Further, it is found that each of the resulting mass operators is equivalent to one which can be derived from an infinite component wave equation. Moreover, it is shown that these mass operators admit spacelike solutions, so that one is confronted with these unphysical solutions in attempting to carry out the program of saturating the current algebra at infinite momentum. However, if these spacelike solutions are uncoupled by the current operator, then one could still saturate the current algebra, but this possibility is closed since it has been shown that the current does in fact connect the spacelike and timelike solutions in all nontrivial cases.

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Erratum

Dynamics of a Double-Peaked Resonance, R. RUSSELL CAMPBELL, PHILLIP W. COULTER, AND GORDON L. SHAW [Phys. Rev. D 2, 1184 (1970)]. The ordinate label η of the top curve of Fig. 1 was accidentally deleted in the printing. In the caption of Fig. 1(a), ‡ should read ρ. In the caption of Fig. 1(b), the equation should read $\delta=\delta_{BW}\{1+0.1/[(s-12)^2+0.5]\}$.