region.

It may be noticed from Eq. (9) that if there were a pure $I=\frac{1}{2}$ resonance with no background, then in its neighborhood the Glauber term would be zero. On the other hand, the new correction would not be zero and would actually correspond to antishadowing.

In principle, one should also consider diagrams such as that of Fig. 1, but with the intermediate pion replaced by a ρ meson.[5] Below a few BeV their contribution is strongly suppressed because the minimum momentum transfer q allowed is

$$q \approx (\mu_\rho^2 - \mu_\pi^2)/2k. \tag{14}$$

According to the eikonal philosophy, for not-too-high energies, this momentum transfer makes the ρ contribution incoherent with the pion over the deuteron radius and, therefore, unimportant. In any case, their inclusion is extremely difficult because it would require knowledge of the phase of the amplitude $A_{\pi N \to \rho N}$. One would then have to resort to such unreliable models as those of Regge or SU(6). It is amusing to note that if one describes the ρ-production amplitude by a one-pion exchange, then the isospin factors are such as to cancel the ρ-contribution completely.

In summary, we have derived a modified Glauber formula which is charge independent, but is as easy to apply as the original. The fact that this modification was not proposed much earlier shows that in such problems a diagrammatic approach may have something to add to an eikonal one.

The author wishes to thank Dr. A. Mueller and Dr. R. F. Peierls for useful discussions and much patience. He is especially indebted to Professor L. Wolfenstein for clarifying many of the ideas. An interesting conversation with Professor R. J. Glauber is also appreciated.

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"REGENERATION" EFFECTS IN ω-φ PRODUCTION*

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According to current ideas,[1] the ω and φ mesons are a coherent mixture of a pure SU(3) octet state φ_0 and a pure SU(3) singlet ω_0. The eigenstates for the freely propagating ω and φ turn out to be approximately[2] $\varphi = (\sqrt{2}\varphi_0 - \omega_0)/\sqrt{3}$ and $\rho = (\varphi_0 + \sqrt{2}\omega_0)/\sqrt{3}$. Now there is no reason to expect these particular linear combinations to be respected in general; for example, in high-energy scattering we might suppose that unitary spin exchange, like isospin exchange, is small, in which case φ_0 and ω_0 are the eigenstates of the scattering. Under these circumstances, the "elastic" scattering amplitudes for φ_0 and ω_0, (if they are not accidentally degenerate), produce "inelastic" reactions of the type $\varphi \to \omega$ by "diffraction dissociation"[3] or, in the terminology of the K^0 mixing theory, by "regeneration" of an ω by an incoming φ.

We would like to point out that effects closely analogous to the K^0 mixing phenomena should, in general, exist for high-energy ω-φ production and to examine cases in which they might be observed.[4] Even though the masses of φ and ω differ by a full 240 MeV, we are justified in using the particle-mixture analogy, since the question of coherence is relative to the phase development of the wave over the scattering system. If we think of an ω or a φ of $K=5$-BeV/c momentum passing through nuclear matter in the lab, the phase difference developed is $\Delta t(\omega_\varphi-\omega_\omega)=\Delta t[(K^2+m_\rho{}^2)^{1/2}-(K^2+m_\omega{}^2)^{1/2}] \simeq \Delta t(m_\varphi{}^2-m_\omega{}^2)/2K=0.22$ rad/F, so that even at this relatively low energy the ω-φ combination can be coherent over a substantial number of nucleons in a nucleus. We also are well justified in neglecting the possible decay in the interaction region since a particle of $M\sim 1$ BeV, $\Gamma\sim 10$ MeV travels $\gamma\Gamma^{-1}=100$ F at 5 BeV.

By using nuclei of different radii, we can vary the "thickness" of our "regenerator" and look for effects as a function of nuclear mass number A. Of course, in a scattering experiment we cannot literally look in the shadow of the nucleus in the way that we observe regeneration of K's behind a macroscopic slab, but the diffracted wave we observe at infinity is essentially constructed from the wave behind the nucleus; so the situation is quite the same.

Nuclear mixing formulation. – The forward photoproduction of the photonlike vector mesons ρ, ω, and φ presents a suggestive situation in which the production reaction itself may be thought of as a species of regeneration due to the absorption of various components virtually present in the photon, and we have shown[5] that such a model gives good agreement with experiments[6] on ρ^0 photoproduction. We can proceed similarly for ω-φ photoproduction (see below). But at the outset, let us take a more conservative point of view and simply note, independently of any specific production mechanism, that since ω and φ have the same quantum numbers [SU(3) aside], the amplitude for $\gamma+p\to\binom{\omega}{\varphi}+p$ can be written as a two-component column vector representing an outgoing linear combination of ω and φ:

$$f_p=\begin{pmatrix} f_{\omega p}\\ f_{\varphi p}\end{pmatrix},$$

$$\frac{d\sigma}{d\Omega}(\gamma+p\to\omega+p)=|f_{\omega p}|^2. \quad (1)$$

(Except when discussing the questions of the relative phase of the waves, we neglect the small kinematic effects due to the difference in momenta at a given energy.)

We anticipate that, like the photo-ρ production, the small angle ω-φ production will be coherent on nuclei. In the usual treatment of high-energy coherent nuclear photoproduction,[7] without mixing complications, we construct the optical model wave function for the outgoing particle of momentum q

$$\Psi^-(x)=\exp(iq\cdot x-D^{\text{out}}/2\lambda), \quad (2)$$

and obtain the coherent nuclear amplitude[7] f_A

$$f_A=kAf_p\left[\frac{1}{\text{vol}}\int_{\text{sphere}}\Psi^-(x)e^{iK\cdot x}d^3x\right],$$
$$\equiv kAf_pQ. \quad (3)$$

The constant k is a scale factor, which we have found by comparison with π- and p-nucleus scattering to be $k^2=1.37$, presumably due to nuclear effects unaccounted for by the impulse approximation. With this adjustment we have found that Eq. (3) gives excellent results.[5] The unusual aspect of the φ-ω problem is that the absorption may mix the φ and ω into each other, and we should treat Eq. (2) as a matrix in the ω-φ space indicated by Eq. (1), with $1/2\lambda$ a matrix with off-diagonal terms, if ω and φ are not the diagonal states in nuclear matter.

The kinematic conditions of high-energy production at small angles on a heavy target give the outgoing mesons the same energy equal to that of the incoming photon, but different momenta of magnitude q_ω and q_φ due to the mass difference $q_\omega-q_\varphi=\Delta_0 M_\omega{}^2-M\ {}^2/2K$. Thus in our stationary wave function corresponding to a definite energy, q is also a matrix, diagonal in the φ-ω basis, with entries q_ω and q_φ.

To find the generalization of Eq. (3), let us consider a wave leaving a uniform slab with a matrix index of refraction. Then, since the eikonal or high-energy approximation amounts to solving a one-dimensional wave equation along a ray, we can simply add up the results for different impact parameters ρ. If the slab is centered on $Z=0$ and has the right-hand edge at $Z=a$, then we want a wave function which becomes an ω or a φ with phase $e^{iqa}\times\binom{1}{0}$ [or $\binom{0}{1}$] at a and satisfies the matrix wave equation as we go back into the slab. Since we deal with a one-dimensional equation, it can be solved by direct exponentiation, just like the time-dependent amplitude used in the K^0 problem, giving

$$\Psi^-(Z)=e^{(iq-1/2\lambda)(Z-a)}e^{iqa}, \quad (4)$$

where separation of the noncommuting exponentials is explicitly necessary.

Now for the nucleus, $a=(R^2-\rho)^{1/2}$ is the distance from $Z=0$ to the edge along the ray, and Eq. (3) becomes

$$f_A = kA\begin{pmatrix}1\\0\end{pmatrix}Qf_p \text{ for } \omega,$$
$$= kA\begin{pmatrix}0\\1\end{pmatrix}Qf_p \text{ for } \varphi, \tag{5}$$

where Q is now a 2×2 matrix, since Ψ^- is given by Eq. (4).

$$Q=\frac{1}{\text{vol}}\int_{\text{sphere}}\exp[-iq(R^2-\rho^2)^{1/2}] \times\exp\{(-iq-1/2\lambda)[Z-(R^2-\rho^2)^{1/2}]\}\exp(+iK\cdot x)d^3x$$

with, in the φ-ω basis[8]

$$q=\begin{pmatrix}q_\omega & 0\\ 0 & q_\varphi\end{pmatrix}, \tag{6}$$

$$1/2\lambda=\langle 1/2\lambda\rangle_{\text{av}}I+(1/2\lambda)^-(\sigma_3\cos 2\theta_N+\sigma_1\sin 2\theta_N),$$

where

$$\langle 1/2\lambda\rangle_{\text{av}}\equiv\tfrac{1}{2}[(1/2\lambda)_{\omega'}+(1/2\lambda)_{\varphi'}],$$

and

$$(1/2\lambda)^-\equiv\tfrac{1}{2}[(1/2\lambda)_{\omega'}-(1/2\lambda)_{\varphi'}].$$

(σ's are the usual Pauli matrices.) $\lambda_{\omega',\varphi'}$ are the effective mean free paths for the combinations $\omega'=\omega\cos\theta_N+\varphi\sin\theta_N$, $\varphi'=-\omega\sin\theta_N+\varphi\times\cos\theta_N$, which are diagonal in nuclear matter, with θ_N the "mixing angle" in nuclear matter.

Diffraction dissociation model.—The diffraction dissociation formalism which we have developed[5] for such reactions gives a definite prediction for f_p in Eq. (1), and allows the whole formulation to be put on a more elegant footing. It consists essentially of pushing the ideas discussed above back onto the production process itself and allowing virtual states "present" in the photon to be "regenerated" by traveling through the target, be it nucleon or nucleus. Thus if we use SU(3), which couples the photon to φ_0, the production amplitude for $\gamma\to\omega,\varphi$ is[4]

$$f=\begin{pmatrix}f_{\omega\omega} & f_{\omega\varphi}\\ f_{\varphi\omega} & f_{\varphi\varphi}\end{pmatrix}\begin{pmatrix}1/m_\omega^{\;2} & 0\\ 0 & 1/m_\varphi^{\;2}\end{pmatrix} \times\begin{pmatrix}(\frac{1}{3})^{1/2}\\ (\frac{2}{3})^{1/2}\end{pmatrix}g_{\gamma\varphi_0}M_{\varphi_0}^{\;2}, \tag{7}$$

where the f matrix is the elastic scattering amplitude for φ-ω on the target, which can be expressed in terms of θ_N and the diagonal amplitudes for φ' and ω'. This more general approach is, in fact, equivalent to Eqs. (3) or (5) when the target is a nucleus. The production amplitude on a nucleon given by Eq. (7) is proportional to the elastic scattering amplitude on the nucleon, and $kAf_{\text{el}}Q$ is the nuclear scattering amplitude in the optical model, so that the photoproduction amplitude does indeed come out proportional to the nuclear elastic-scattering amplitude.

Now, if we assume that ω-φ high-energy scattering is diffractionlike and predominantly imaginary at small angles, then the f matrix in Eq. (7) is proportional by the optical theorem to a total cross-section matrix $f=(k/4\pi)\sigma_{\text{tot}}$, and we anticipate that the components of f_p in Eq. (1) are both imaginary; furthermore, with our sign for the mixing angle $\varphi_0=(\sqrt{2}\varphi+\omega)/\sqrt{3}$, the relative sign[2] is plus.

The data indicate that φ production on protons is small, relative to ω production; this we have interpreted, in terms of Eq. (7), as essentially indicating the $\sigma_{\text{tot}}(\varphi'p)$ is small [for $\theta_N=0$ we estimated that at $K<5$ BeV, $\sigma_{\text{tot}}(\omega p)\sim 65$ mb, $\sigma_{\text{tot}}(\omega p)\leqslant 25$ mb] and that θ_N cannot be too near the actual φ-ω mixing angle, or we cannot easily suppress the large φ component in φ_0. Qualitatively, the dissociation model implies that high-energy amplitudes for $\gamma+p\to p+\omega,\varphi,\rho$ behave like elastic diffraction scattering amplitudes. In what follows we shall use the experimental evidence that the ω component of f_p in Eq. (1) is substantially greater than the φ component, and the dissociation predictions that the amplitudes are predominantly relatively real and positive and that $\sigma_{\text{tot}}(\omega'p)$ is substantially bigger than $\sigma_{\text{tot}}(\varphi'p)$.

Effects in ω-φ production.—The matrix Q obviously contains much interesting structure (including some possible small φ-ω variations with production angle), but let us consider forward production; two cases may be treated directly.

(A) Nuclear-scattering diagonal for the physical φ and ω, $\theta_N=0$. (This is effectively the case if the φ'-p cross sections are approximately equal.) In this case the matrix formulation is an unneccessary formality since the φ and ω are decoupled and the simple Eq. (3) applies. However, even here it should be realized that at the lower energies there will be some (smooth) variation in the forward cross section versus A because of the finiteness of the momentum transfer, which cuts down the coherent cross section as we increase the nuclear radius. In Fig. 1(a) we show φ production per nucleon for this case at 5 BeV ($\Delta_0 = 104$ MeV) for various values of the total φ-p cross section. There is essentially a (longitudinal) diffraction "wiggle" except that, contrary to general practice, we obtain it by varying the nuclear size instead of the momentum transfer. By 15 BeV, however, the curves flatten out and simply rise smoothly throughout the periodic table. It is worth noting that the smallness of $\sigma(\varphi p)$ found in the dissociation model immediately implies a rise in φ/ω production ratio for $\theta_N = 0$, since the φ's will have a small absorption. This effect is peculiar to the dissociation concept and does not necessarily follow generally. In fact, if the cross sections were large enough so that the elastic f's in Eq. (7) saturate at geometric values on heavy nuclei, then the dissociation model predicts that the φ/ω ratio would just approach their ratio in the photon, $2(M_\omega/M_\varphi)$.[4] Unfortunately, for the cross sections we have in mind, the rise towards saturation is slow At 15 BeV, for $\sigma_{\text{tot}}(\omega p) \approx 50$ mb and $\sigma_{\text{tot}}(\varphi p) \approx 26$ mb, the φ/ω ratio increases by a factor of 2 from H to Pb, by a factor of 4 for $\sigma_{\text{tot}}(\varphi p) \approx 13$ mb.

(B) High energy, so that the ω-φ mass difference is unimportant and q is effectively proportional to the unit matrix $(q_\omega - q_\varphi)R < 1$ even on the largest nuclei. In this limit we can simply find the "form factors" Q for the φ' and ω' and then rotate back to the physical states, so that in the φ-ω basis

$$Q = \tfrac{1}{2}(Q_{\omega'} + Q_{\varphi'})I + \tfrac{1}{2}(Q_{\omega'} - Q_{\varphi'}) \times (\sigma_3 \cos 2\theta_N + \sigma_1 \sin 2\theta_N). \qquad (8)$$

In general, however, we cannot neglect the difference in phase development between the

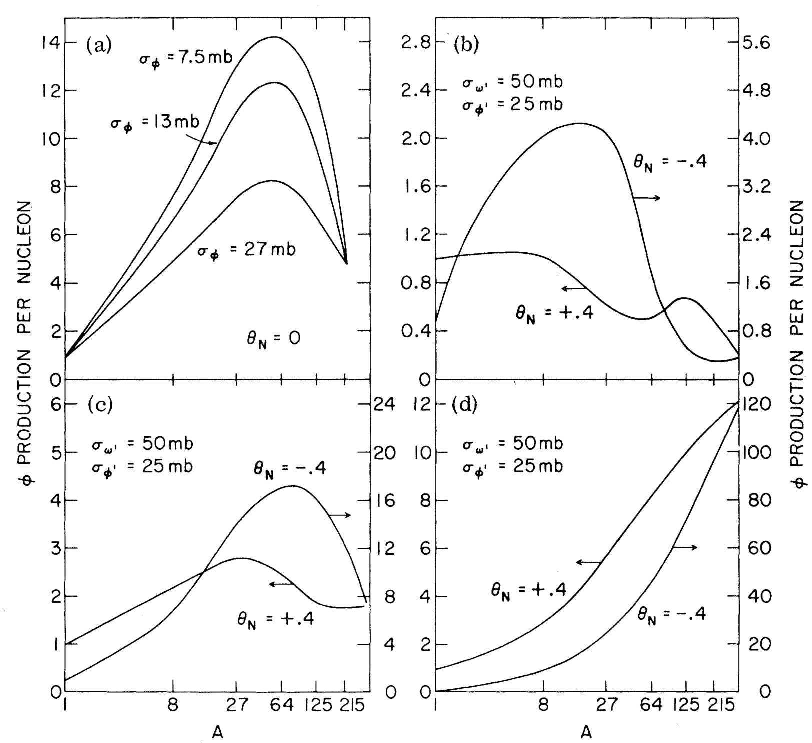

FIG. 1. Forward φ production per nucleon versus nuclear mass number A for various values of the parameters. (a) No mixing, at 5 BeV. The cross section per nucleon drops for higher mass numbers, where we have $\Delta R > 1$. (b)-(d) A moderate mixing angle, with both constructive and destructive interference, for 3, 5, and 15 BeV, respectively. The structure smooths out as we increase the energy. Note that in the 15-BeV constructive case, for small mass number, the cross section is about equal to the fully coherent value of A per nucleon, despite the absorption.

components of definite mass (just as the K_1-K_2 mass difference can lead to effects such as oscillations in the strangeness) and the matrix exponential in Eq. (6) must be expanded; giving for forward production

$$Q=\frac{2\pi}{\mathrm{vol}}\int\rho d\rho\,\exp[i\Delta(R^2-\rho^2)^{1/2}(i\Delta+1/2\lambda)^{-1}\times\{1-\exp[2(i\Delta+1/2\lambda)(R^2-\rho^2)^{1/2}]\},$$

where

$$\Delta=\begin{pmatrix}k-q_\omega & 0\\ 0 & k-q_\varphi\end{pmatrix}$$

and k = photon momentum.

If we write

$$\Delta=\bar{\Delta}+\Delta^-\sigma_3,$$

$$1/2\lambda=\langle 1/2\lambda\rangle_{\mathrm{av}}+(1/2\lambda)^-(\sigma_3\cos2\theta_N+\sigma_1\sin2\theta_N),$$

and

$$i\Delta+1/2\lambda=\bar{\Lambda}+\vec{\Lambda}\cdot\vec{\sigma},$$

then

$$\bar{\Lambda}=i\bar{\Delta}+\langle 1/2\lambda\rangle_{\mathrm{av}},\quad \Lambda_3^-=i\Delta^-+(1/2\lambda)^-\cos2\theta_N,$$

and

$$\Lambda_1=(1/2\lambda)^-\sin2\theta_N.$$

Now define $\Lambda^-=(\Lambda_1^2+\Lambda_2^2)^{1/2}$, and then the exponential can be expanded:

$$\exp[-2(i\Delta+1/2\lambda)(R^2-\rho^2)^{1/2}]=\tfrac{1}{2}\{\exp[-2(\bar{\Lambda}+\Lambda^-)(R^2-\rho^2)^{1/2}]+\exp[-2(\bar{\Lambda}-\Lambda^-)(R^2-\rho^2)^{1/2}]$$
$$+(\vec{\Lambda}\cdot\vec{\sigma}/\Lambda^-)\exp[-2(\bar{\Lambda}+\Lambda^-)(R^2-\rho^2)^{1/2}]-\exp[-2(\bar{\Lambda}-\Lambda^-)(R^2-\rho^2)^{1/2}]\},$$

so that the integral may be evaluated to give the general result for forward production, including both the effects of mixing due to the off-diagonal absorption and the φ-ω mass difference.

Interesting effects may be observable in φ production. Since "direct" φ production is small, a substantial part of the final φ amplitude will be due to ω's converted into φ's so that the direct and regenerated amplitudes may interfere. Since φ production on protons is $\lesssim\frac{1}{10}$ ω production, let us take in our examples $f_p \propto \binom{(10)^{1/2}}{+1}$. (We should emphasize, however, that in our diffraction model, f_p is actually closely related to the parameters determining the behavior of the nuclear cross sections and so, experimental consistency checks are possible.) In Fig. 1(b) we show the φ production per nucleon at 3 BeV for $\theta_N=0.4$, $\sigma_{\omega'}=50$ mb, $\sigma_{\varphi'}=25$ mb, right or left scales applying according to whether the regenerated amplitude adds constructively or destructively. At this relatively low energy where $\Delta_\varphi=173$ MeV and $\Delta_\omega=102$ MeV, the difference in the φ-ω phase development has an important effect as we see by comparing with 5 BeV [Fig. 1(c)]. In the 15-BeV constructive-interference case, we see that the φ production becomes as great as the fully coherent value of A per nucleon on the light nuclei. This is an indication of an amusing possibility of a supercoherent effect in which the regenerated amplitude tends to add to the direct amplitude, with the result that the cross section actually exceeds the fully coherent value. To take an extreme case, let θ_N be -0.67, the "vacuum" mixing value, $\sigma_{\omega'}=65$ mb, and $\sigma_{\varphi'}=7$ mb. Figure 2(a) shows the results at 3, 5, and 15 BeV, which rise to the huge value of more than twice the fully coherent value at 15 BeV. [On the other hand, with such a large value of θ_N, such widely disparate values of the σ's are not necessary for structure as the curves in Fig. 2(b) with $\theta_N=\frac{2}{3}$, destructive interference, indicate.] Note that since we believe $\sigma(\varphi'p)<\sigma(\omega'p)$, we have $(Q_{\omega'}-Q_{\varphi'})<0$ in Eq. (8); thus, with our definite choice of phases, constructive interference indicates $\theta_N<0$ (for θ_N not very large), or that φ' tends to be like φ_0.

While we cannot, because of the very implicit combination of effects of the various unknown parameters in the calculation, give a general description of what may be expected experimentally, the examples indicate that it may be quite possible to see interesting effects in φ production without very extreme values of the parameters. As for the idealizations that the optical potential is purely imaginary (i.e., absorptive) and that the nucleus has a sharp edge, a real part for the forward scattering amplitude on protons for the outgoing particle of the general order found in high-energy elastic-scattering experiments will lead to a local momentum for

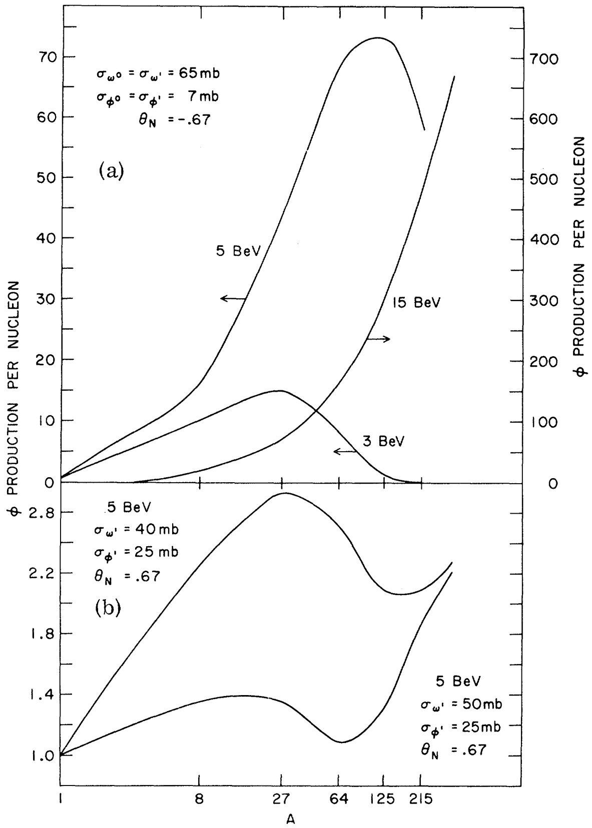

FIG. 2. Forward φ production per nucleon for large mixing, $|\theta_N| = 0.67$. (a) With $\theta_N = -0.67$, about roughly the ω-φ mixing angle, and the extreme case of widely differing cross sections, we have the possibility of a large supercoherent effect. (b) Structure with a large mixing angle and relatively small differences in the cross sections.

the particle in the nucleus differing from its vacuum value by several MeV[7]; thus, this is not a problem until very high energies where the momentum transfer can get this small. The effect of some smoothness in the nuclear edge may be simulated by assigning some uncertainty to the abscissa in our graphs; since the oscillations are usually quite broad, it does not appear that this will wipe them out. We wish, then, to emphasize that the experimental study of coherent photoproduction of φ's through the periodic table and at a variety of energies is of the greatest interest.

If the physical ω and φ are not the diagonal states in nuclear matter, it is clear that there also should be some effects in ω-φ production by incident π's or k's or in $\bar{p}$ annihilation. Since, however, these reactions (involving a charge or baryon exchange) cannot be coherent, the phenomenon simply amounts to the observation that there is some $\omega \rightleftarrows \varphi$ conversion in the rescattering of the outgoing particle. In any case, if we take a reaction such as $\pi + N \to \omega + N'$ (where N' could be a nucleon or an N^*), which seems to be producing predominantly ω's and few φ's, then we should expect an increase in the relative number of φ's by using nuclear targets, if indeed φ and ω are somewhat mixed in nuclear matter.

The observation of such an effect would be a simple qualitative indication that there is substantial mixing in nuclear matter and thus immediately allow us to draw conclusions about the nature of high-energy ω-φ nucleon scattering. Ideas similar to those given here in this paper also apply, of course, to other pairs of particles with identical or almost identical quantum numbers. We would like to thank J. Yellin and R. Serber for interesting discussions.

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[8]We use a uniform-sphere picture of the nucleus with cylindrical coordinates. Z is the direction of outgoing particle, ρ = impact parameter, R = nuclear radius $= A^{1/3} \times 1.2$ F, $D^{\mathrm{out}} = (R^2 - \rho^2)^{1/2} - Z$, and K = momentum of incident photon. It is assumed that the forward scattering amplitude for the outgoing particle is predominantly imaginary, and spin and isospin independent, in which case the absorption is characterized by a mean free path $\lambda^{-1} = \rho\sigma_{\mathrm{tot}}$, where ρ = density of nucleons and σ_{tot} = total cross section on a nucleon.