

Mass dependence of inclusive nuclear ϕ photoproduction

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Abstract

Based on a prior determination of the ϕ selfenergy in a nuclear medium we perform a theoretical study of inclusive ϕ photoproduction in nuclei, looking at the A dependence of the cross sections for different ϕ momenta. We find sizeable reductions in the nuclear cross sections with respect to the elementary one, using a ϕ selfenergy which implies a width about six times the free one at normal nuclear density. The calculations are done to match the set up for an ongoing experiment at *SPring8/Osaka* which should provide valuable information on the renormalization of the ϕ properties in nuclei.

1 Introduction

The renormalization of the properties of the vector mesons in nuclear matter has received much attention, particularly the ρ meson, (see [1] for a review). The ϕ meson has received comparatively less attention, but it turns out to be a well suited tool to understand the dynamics of the vector mesons in matter since the changes of some of the ϕ properties are comparatively larger than those of the ρ meson, what would make its experimental observation in principle easier. The change of the ϕ properties in the medium at finite temperature and/or density has been studied in several approaches like effective Lagrangians [2, 7, 4, 5, 6, 3] and QCD sum rules [8, 9, 10, 11]. Particularly, the ϕ width modification in matter has been subject of study in dropping meson mass scenarios [6, 12, 13, 14, 15, 16], as a result of collisional broadening through ϕ -baryon [17] or

ϕ -meson [18, 19] scattering processes and by modifying the $\bar{K}K$ decay channel by means of in-medium kaon selfenergies [7, 20]. The majority of these works point at a sizeable renormalization of the ϕ width and a small mass shift.

Indeed, in [3, 7] it was found that the ϕ width at normal nuclear matter density, ρ_0 , was about one order of magnitude larger than the free width. Particularly, the results of [7] were based on the use of a K^- selfenergy in nuclear matter which accounted for Pauli blocking correction in the intermediate K^-N states in the scattering equation. Further studies considering the selfconsistency of the $\bar{K}N$ in-medium scattering equation were done in [21, 22] to determine the K^- selfenergy. Following those results a reanalysis of the ϕ width in the medium was done in [20], resulting in values around 22 MeV at $\rho = \rho_0$. A more recent evaluation, using the K^- selfenergy of [22] and improving on some approximations of [20] was done in [23], resulting in a width of about 28 MeV at $\rho = \rho_0$ and a small shift in the mass of about -6 MeV.

Several proposals to test the ϕ properties in a nuclear medium have been done, for instance, regarding the detection of its decay products in $A-A$ and $p-A$ collisions [24, 11]. Particularly a method to analyse the invariant mass spectra of the emitted charged kaon pairs was studied in [25]. Having into account that the calculations in [3, 7, 20, 23] were done for a ϕ at rest in the nucleus, it has been suggested to determine this large width using reactions like $\pi^-p \rightarrow \phi n$ in nuclei [7] and ϕ photoproduction in nuclei [26], assuming the ϕ to emerge with a small momentum. The latter proposal is possible by means of the elementary $\gamma N \rightarrow \phi N$ reaction with the ϕ going in the backward direction and the help of the Fermi motion of the nucleons. The small cross section for these particular events makes this experiment difficult and in addition it was shown in [27] that the distribution of the K^+K^- pairs from the Coulomb interaction in heavy nuclei removed the changes in the ϕ width in the K^+K^- invariant mass that one might expect from the changes of the ϕ width in the medium. Another possibility, not previously considered, to investigate the changes of the ϕ width is to look for ϕ loss of flux in ϕ nuclear photoproduction. The A dependence in this loss is related to the ϕ width in the medium. The drawback is that in this experiment the largest fraction of the ϕ comes out with a momentum of the order of 1500 MeV/c in the ongoing experiment at *SPring8/Osaka* [28]. The combination of experiment and theoretical models, which can make the extrapolation of ϕ at rest to ϕ with finite momenta, can thus make the experiment useful to learn about ϕ nuclear properties.

With this aim, we present here a theoretical calculation of the A dependence of the cross section for ϕ photoproduction in nuclei which could test the models of [3, 7, 20, 23]. Our work will be based on the models for the ϕ selfenergy of Oset-Ramos (OR) [20] and Cabrera-Vicente (CV) [23], after a proper extension to account for the finite momentum.

2 ϕ photoproduction cross section

Let $\Pi(p, \rho)$ be the ϕ selfenergy in a nuclear medium as a function of its momentum, p , and the nuclear density, ρ . We have

$$\frac{\Pi}{2\omega} \equiv V_{\text{opt}} = \text{Re}V_{\text{opt}} + i \text{Im}V_{\text{opt}} , \quad (1)$$

and hence

$$\frac{\Gamma}{2} = -\text{Im}\frac{\Pi}{2\omega} \quad ; \quad \Gamma = -\frac{\text{Im}\Pi}{\omega} \equiv \frac{dP}{dt} , \quad (2)$$

where ω is the ϕ energy and P is the ϕ decay probability, including nuclear quasielastic and absorption channels. Hence, we have for the probability of loss of flux per unit length

$$\frac{dP}{dl} = \frac{dP}{v dt} = \frac{dP}{\frac{p}{\omega} dt} = -\frac{\text{Im}\Pi}{p} . \quad (3)$$

The nuclear cross section for inclusive ϕ photoproduction will then be

$$\frac{d\sigma_A}{d\Omega} = \int d^3\vec{r}\rho(r) \frac{d\sigma}{d\Omega} e^{-\int_0^\infty dl \frac{1}{p} \text{Im}\Pi(p,\rho(r'))} \quad (4)$$

where $\frac{d\sigma}{d\Omega}$ and $\frac{d\sigma_A}{d\Omega}$ are the elementary and nuclear differential cross section and $\vec{r}' = \vec{r} + l\frac{\vec{p}}{|\vec{p}|}$ with \vec{r} the ϕ production point inside the nucleus. The exponential factor in Eq. (4) represents the survival probability of the ϕ meson in its way out of the nucleus. If the ϕ did not get absorbed inside the nucleus then we would get the typical result for an electromagnetic reaction $\frac{d\sigma_A}{d\Omega} = A\frac{d\sigma}{d\Omega}$. Eq. (4) relies in the eikonal approximation which is accurate for the large ϕ momentum involved in the process. In practice one might expect small corrections to the formula, even if the ϕ did not decay, from two sources:

- i) Distortion of the ϕ trajectory because of the real part of the potential
- ii) Change of direction and energy of the ϕ in quasielastic collisions $\phi N \rightarrow \phi' N'$.

The first effect should be negligible since the real part of the potential is so weak that only modifies the mass of the ϕ in about 6 MeV [3, 23]. The second effect should be very weak too, account taken of the lack of direct coupling ϕNN because of the OZI rule. On the other hand, the effect from this source would simply lead to a slight change in the ϕ direction but not the disappearance of the ϕ . This simply means that collecting the ϕ in a narrow cone along the forward direction, where practically all the ϕ go, both because of the kinematics of the lab variables and the extremely forward direction of the $\gamma N \rightarrow \phi N$ cross section [29], would guarantee that the ϕ undergoing quasielastic collisions in the nucleus are accounted for in the nuclear cross section. In order to adjust to this experiment we should not remove theoretically the events in which there are ϕ quasielastic collisions. This will be done by including in $\text{Im}\Pi$ only the ϕ absorption events.

The integral in Eq. (4) does not depend on the direction of the ϕ momentum and thus we have

$$P_{\text{out}} \equiv \frac{\sigma_A}{A\sigma} = \frac{1}{A} \int d^3\vec{r}\rho(r) e^{\int_0^\infty dl \frac{1}{p} \text{Im}\Pi(p,\rho(r'))} \quad (5)$$

and this is the magnitude which we would evaluate as a function of p and A , which can be interpreted as the probability for a ϕ to go out of the nucleus. The density profiles, $\rho(r)$, for the different nuclei have been taken from [30].

3 Evaluation of the ϕ selfenergy at finite momentum

For the evaluation of $\text{Im}\Pi$ in nuclear matter we are going to use the OR and CV models which are based on chiral $SU(3)$ dynamics considering the K and \bar{K} in-medium properties. In these works the evaluations were done for a ϕ at rest. Since in the present work we are dealing with ϕ mesons with momentum up to 2 GeV, an extension of the models to finite momenta is mandatory. In the following subsections we summarize these two models stressing the modifications to consider the finite ϕ momentum.

3.1 Extension of the Oset-Ramos model to finite momentum

The OR accounts for $\phi \rightarrow K^+K^-$ and $\phi \rightarrow K^0\bar{K}^0$ decay diagrams where the kaons are renormalized in the medium due to s-wave and p-wave interactions, Figs. 1 and 2 respectively.

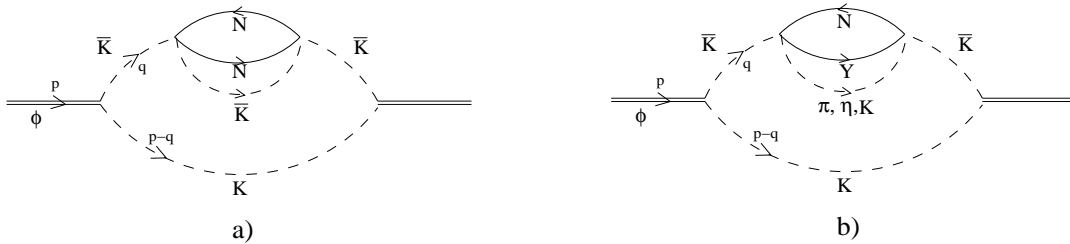


Figure 1: Diagrams contributing to the ϕ selfenergy coming from the s-wave $\bar{K}N$ interaction. Y represents Λ , Σ or Ξ .

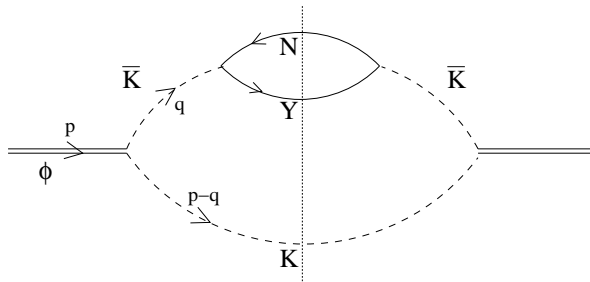


Figure 2: Diagrams contributing to the ϕ selfenergy coming from the p-wave $\bar{K}N$ interaction. Y represents Λ , Σ or $\Sigma^*(1385)$.

To the diagrams of Fig. 2 one must add the vertex corrections shown in Fig. 3 for consistency with gauge invariance.

In the evaluation of the previous diagrams, the K^+ , K^0 , selfenergy is taken as in [31] following the $t\rho$ approximation

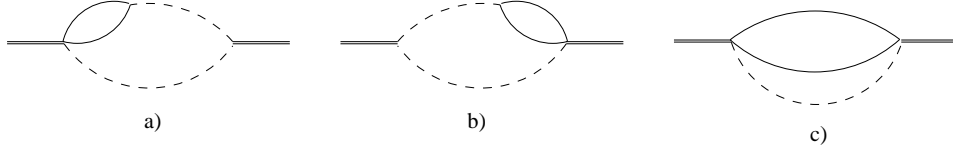


Figure 3: Diagrams accounting for vertex corrections.

$$\Pi_{K^+,K^0} = 0.13 m_K^2 \frac{\rho}{\rho_0} \text{ MeV} . \quad (6)$$

Concerning the \bar{K} , a large contribution to the ϕ selfenergy comes from the p-wave and vertex correction diagrams. On the other hand, since the K^- p-wave selfenergy is proportional to \vec{q}^2 , it is likely to be more affected inside the loops by the finiteness of the ϕ momentum. Hence we evaluate explicitly the ϕ selfenergy from the p-wave K^- selfenergy diagrams at finite ϕ momentum and take the same part coming from the K^- s-wave selfenergy as obtained for zero momentum. This latter part produces a contribution to the imaginary part of the ϕ selfenergy that can be parametrized as

$$\text{Im}\Pi^{(s)} = -7.62\rho/\rho_0 \omega(p) \text{ MeV} \quad (7)$$

We now evaluate the contribution to $\text{Im}\Pi(p, \rho)$ of the \bar{K} p-wave selfenergy diagrams, depicted in Fig. 2.

In Fig. 2, K represents K^+ or K^0 , \bar{K} represents K^- or \bar{K}^0 and Y the hyperons Λ , Σ and $\Sigma^*(1385)$. The ϕ selfenergy coming from these diagrams at finite momentum can be obtained as

$$\begin{aligned} -i\Pi^{(p)}(p, \rho) &= 2 \int \frac{d^4q}{(2\pi)^4} iD(q)iD(q)(-i)\Pi_{\bar{K}}^{(p)}(q, \rho)i\tilde{D}(p-q, \rho) \overline{\sum} V_{\phi K\bar{K}} V_{\phi K\bar{K}} = \\ &= 2g_\phi^2 \int \frac{d^4q}{(2\pi)^4} D(q)^2 \Pi_{\bar{K}}^{(p)}(q, \rho) \tilde{D}(p-q, \rho) \frac{4}{3} \left[\frac{(q \cdot p)^2}{m_\phi^2} - q^2 \right] \end{aligned} \quad (8)$$

where $D(q) = (q^2 - m_K^2 + i\epsilon)^{-1}$ is the K^- , \bar{K}^0 , propagator, $\tilde{D}(p-q) = ((p-q)^2 - m_K^2 - \Pi_{K^+})^{-1}$ is the K^+ , K^0 , propagator with the selfenergy of Eq. (6). In the evaluation of Eq. (8) we have used that $V_{\phi K\bar{K}} = -ig_\phi \epsilon_\mu (p_K^\mu - p_{\bar{K}}^\mu)$, with $g_\phi = 4.57$. The factor 2 at the beginning of the equation is due to the possibility of having K^-K^+ or \bar{K}^0K^0 in the intermediate states.

To obtain $\text{Im}\Pi$ we can apply the following Cutkosky rules for the cut represented by the dotted line of Fig. 2

$$\begin{aligned} \Pi(k) &\rightarrow 2i \text{Im}\Pi(k) \\ D(k) &\rightarrow 2i \Theta(k^0) \text{Im}D(k) \end{aligned}$$

which leads, after performing the q^0 integration, to

$$\text{Im}\Pi^{(p)}(p, \rho) = g_\phi^2 \int \frac{d^3\vec{q}}{(2\pi)^3} \left(\frac{1}{q^{0^2} - \vec{q}^2 - m_K^2} \right)^2 \frac{1}{\tilde{\omega}_{p-q}} \text{Im}\Pi_{\bar{K}}^{(p)}(q, \rho) \frac{4}{3} \left[\frac{(q \cdot p)^2}{m_\phi^2} - q^2 \right] \Theta(q^0) \quad (9)$$

with $\tilde{\omega}_{p-q} = \sqrt{(\vec{p} - \vec{q})^2 + m_K^2 + 0.13m_K^2 \frac{\rho}{\rho_0}}$, $q^0 = p^0 - \tilde{\omega}_{p-q}$ and Θ is the step function. The \bar{K} p-wave selfenergy, $\text{Im}\Pi_{\bar{K}}^{(p)}(q)$, is evaluated in [22] and is a function of the $\bar{K}NY$ vertices and the Lindhard functions for hyperons. In Eq. (9) we also add, as in Ref. [7], a form factor for each kaon-baryon vertex of dipole type, $[\Lambda^2/(\Lambda^2 - q^2)]^2$, with $\Lambda = 1.05$ GeV.

Finally, the evaluation of the vertex corrections of Fig. 3 is done in an analogous way. In this case we need the ϕKNY vertex functions given by

$$\begin{aligned} V_{\phi KNY} &= g_\phi \tilde{V}_{\bar{K}NY} \vec{\sigma} \cdot \vec{\epsilon}^{(\phi)}; & Y = \Lambda, \Sigma \\ V_{\phi KN\Sigma^*} &= g_\phi \tilde{V}_{\bar{K}N\Sigma^*} \vec{S}^\dagger \cdot \vec{\epsilon}^{(\phi)}, \end{aligned} \quad (10)$$

where $\tilde{V}_{\bar{K}NY}$ and $\tilde{V}_{\bar{K}N\Sigma^*}$ are coefficients given in [20]. The resulting expression for the imaginary part of the ϕ selfenergy coming from the vertex corrections is the same as Eq. (9) but doing the following substitution

$$D(q)^2 \vec{q}^2 \frac{4}{3} \left[\frac{(q \cdot p)^2}{m_\phi^2} - q^2 \right] \longrightarrow D(q) \frac{4}{3} \left[\vec{q}^2 - \frac{(p \cdot q)(\vec{p} \cdot \vec{q})}{m_\phi^2} \right] + \left(1 + \frac{\vec{q}^2}{3m_\phi^2} \right). \quad (11)$$

3.2 Extension of the Cabrera-Vicente model to finite momentum

The ϕ selfenergy from $\bar{K}K$ loop diagrams, calculated in [7, 23] for a ϕ at rest in a nuclear medium, is given by

$$\Pi_{\phi}^{\bar{K}K}(p^0, \vec{p}; \rho) = 2ig_\phi^2 \frac{4}{3} \int \frac{d^4q}{(2\pi)^4} \left[\frac{(p \cdot q)^2}{m_\phi^2} - q^2 \right] D_K(p^0 - q^0, \vec{p} - \vec{q}; \rho) D_{\bar{K}}(q^0, \vec{q}; \rho) \quad (12)$$

for a ϕ meson with a momentum p .

A spectral representation [7] which sums up to all orders the insertion of irreducible kaon selfenergy terms is used in the kaon propagator. The imaginary part of $\Pi_{\phi}^{\bar{K}K}$, which is of interest in the present work, can be written as

$$\begin{aligned} \text{Im}\Pi_{\phi}^{\bar{K}K}(p^0, \vec{p}; \rho) &= -\frac{1}{4\pi} g_\phi^2 \frac{4}{3} \int_0^\infty dq \vec{q}^2 \int_{-1}^1 du \frac{1}{\tilde{\omega}_{p-q}} \\ &\quad \left[\frac{(p \cdot q)^2}{m_\phi^2} - q^2 \right]_{q^0=p^0-\tilde{\omega}_{p-q}} S_{\bar{K}}(p^0 - \tilde{\omega}_{p-q}, \vec{q}; \rho) \Theta(p^0 - \tilde{\omega}_{p-q}), \end{aligned} \quad (13)$$

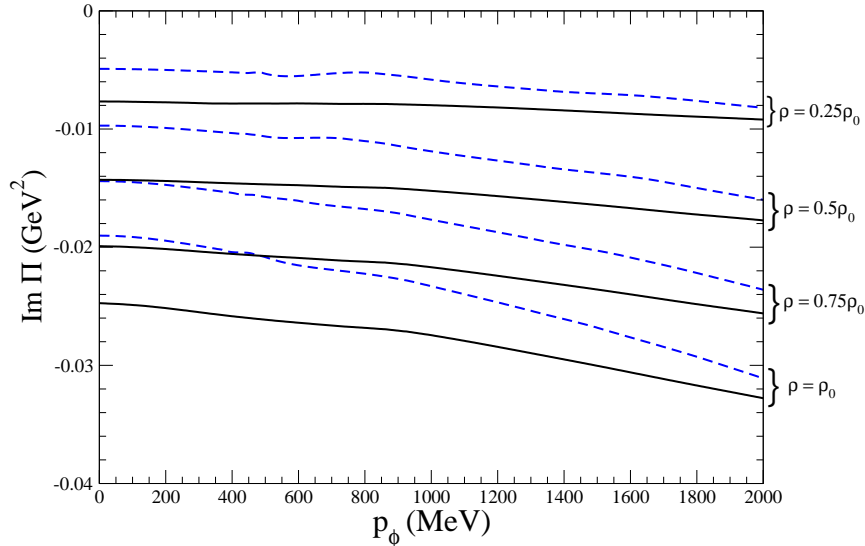


Figure 4: Imaginary part of the in medium ϕ selfenergy, without the inclusion of the free part, as a function of the momentum of the ϕ and the density, for the finite momentum modification of model [20] (dashed line) and [23] (solid line).

where $u = \vec{p} \cdot \vec{q} / |\vec{p}| |\vec{q}|$ and $S_{\bar{K}(K)}$ is the spectral function of the $\bar{K}(K)$ meson,

$$S_{\bar{K}(K)}(q^0, \vec{q}; \rho) = -\frac{1}{\pi} \frac{\text{Im} \Pi_{\bar{K}(K)}(q^0, \vec{q}; \rho)}{|(q^0)^2 - \vec{q}^2 - m_K^2 - \Pi_{\bar{K}(K)}(q^0, \vec{q}; \rho)|^2}. \quad (14)$$

The main differences with respect to the approach described in section 3.1 are the following:

- Eq. (13) considers higher order effects in the insertion of the kaon selfenergies as compared to Eqs. (7) and (9). Moreover, the calculation at finite momentum is considered for both the s- and p-wave \bar{K} selfenergy contributions. Note also that Eq. (13) includes the ϕ free width into the $\bar{K}K$ channel, which will be subtracted to keep only the nuclear effects in the ϕ selfenergy.

- We use the p-wave \bar{K} selfenergy given in [23] which, starting from the result in [20], takes into account an improvement of the relativistic recoil corrections of the $\bar{K}NY$ vertices. In addition, these vertices carry a static form factor of the form $[\Lambda^2 / (\Lambda^2 + \vec{q}^2)]^2$, with $\Lambda = 1.05$ GeV.

- The vertex correction diagrams considered in section 3.1 are also calculated here at finite ϕ momentum, with the prescription of using the fully dressed kaon propagators in the medium.

4 Results and discussion

In Fig. 4 we show the imaginary part of the ϕ selfenergy as a function of the ϕ momentum for different nuclear densities and for the two models described in section 3.1 (dashed line)

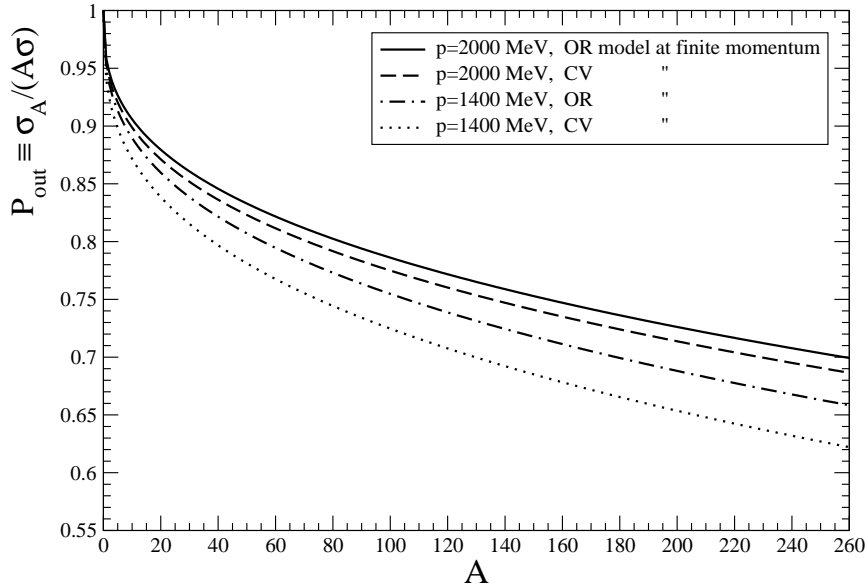


Figure 5: $\sigma_A/(A\sigma)$ as a function of the nuclear mass number for two different momentum of the ϕ calculated with the two models described in the text.

and 3.2 (solid line). The contribution to $\text{Im}\Pi$ coming from the free ϕ decay into $K\bar{K}$, non density dependent, has been subtracted from the full $\text{Im}\Pi$ since the $K\bar{K}$ coming from the free decay would be detected and counted as a ϕ event, hence it does not contribute to the loss of flux required in the argument of the exponential in Eq. (5). Despite the visible differences in Fig. 4 between the two models at zero momentum, the trend of the plot as the momentum increases is very similar. This initial discrepancy at zero momentum is expectable due to the differences in the treatment of the s-wave contribution, the relativistic recoil corrections and the fact that the model of section 3.2 goes beyond the first order in density, implicitly considered in the model of section 3.1. The differences between the two models in Fig. 4, however, are indicative of the intrinsic theoretical uncertainties.

In Fig. 5 we show the ratio $\sigma_A/(A\sigma)$, which represents the probability of one photo-produced ϕ to go out the nucleus, calculated theoretically as a function of A and for two different momenta. Again, a comparison between the two models is shown. We observe that for $p = 1400$ MeV and heavy nuclei this ratio can be of the order of 0.65, indicating a clear ϕ loss of flux which should be identifiable experimentally.

In Fig. 6 we show instead the results for $\sigma_A/(A\sigma)$ as a function of the ϕ momentum for two nuclei and comparing the two theoretical models. In this figure we observe that the amount of ϕ flux lost is larger for smaller ϕ momentum. This is logical to expect since the probability of ϕ decay per unit length is $-\text{Im}\Pi/p$, which is larger for small momenta because of the factor $1/p$. This occurs in spite of the fact that $-\text{Im}\Pi$ decreases with p , as can be seen in Fig. 4, because this p dependence is weaker than that of $1/p$. The set up of the experiment at *SPring8/Osaka* [28] is such that it produces ϕ momenta around 1500 MeV, using photons from 1.4 to 2.4 GeV, coming from ϕ production in the forward direction

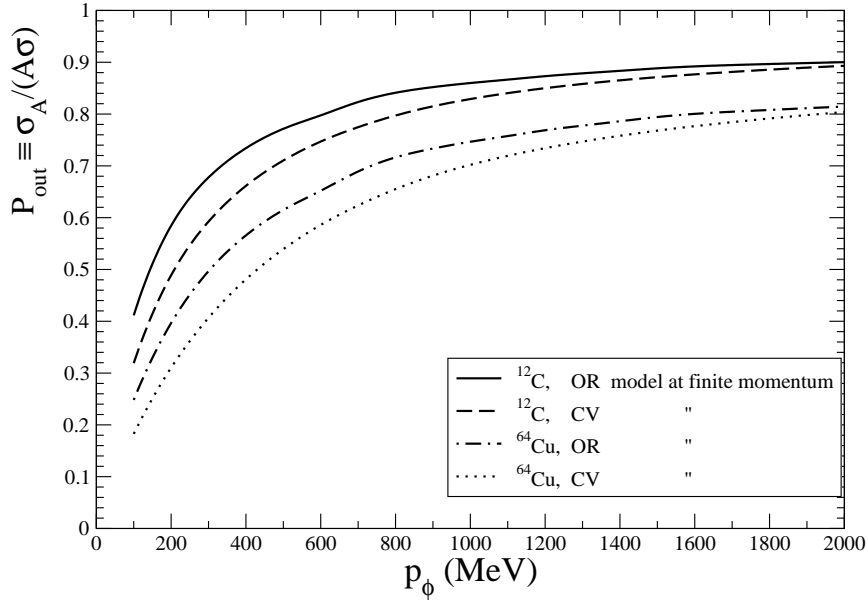


Figure 6: $\sigma_A/(A\sigma)$ as a function of the momentum of the ϕ for two different nuclei calculated with the two models described in the text.

in the CM frame. A possibility to extract further information is to use photons of lower energy, around, 1.6 GeV, not far from threshold which leads to ϕ of around 100 MeV in the CM frame and around 1000 MeV in the lab frame. As one can see in Fig. 6, in a nucleus like ^{64}Cu the depletion factor goes from 0.8 at $p = 2000$ MeV to 0.7 at $p = 1000$ MeV. This latter factor would be 0.57 for $A \simeq 240$.

So far the only nuclear effects considered have been the ones related to the absorption of the ϕ meson. At this point it is worth mentioning that other nuclear effects regarding the production mechanism may lead to a further ϕ loss of flux or change in the ϕ distribution. These other nuclear effects are mainly the Pauli blocking of the final nucleon and the Fermi motion of the initial one. The first one may lead to a reduction of the ϕ flux in comparison to the free case because a certain amount of events are forbidden due to the Pauli blocking of the final nucleon. On the other hand the initial Fermi motion can distort the distribution of the final ϕ mesons. In order to estimate the possible flux reduction due to these sources, we have included in the integrand of Eq. (5) a factor $G(Q, \rho)$ which considers a Fermi average of these effects [32]:

$$G(Q, \rho) = 1 - \Theta(2 - \tilde{Q}) \left(1 - \frac{3}{4}\tilde{Q} + \frac{1}{16}\tilde{Q}^3 \right) \quad (15)$$

where $\tilde{Q} = |\vec{Q}|/k_F$ with Q the momentum transfer of the nucleon and $k_F = (\frac{3}{2}\pi^2\rho(r))^{1/3}$ is the Fermi momentum of the nucleons.

In Fig. 7 we show $\sigma_A/(A\sigma)$, with the CV model, as a function of the ϕ momentum for ^{64}Cu . In solid line we show the result without considering the effect of Eq. (15), (i.e., the same as in Fig. 6). The other lines represent the results considering the Pauli effect

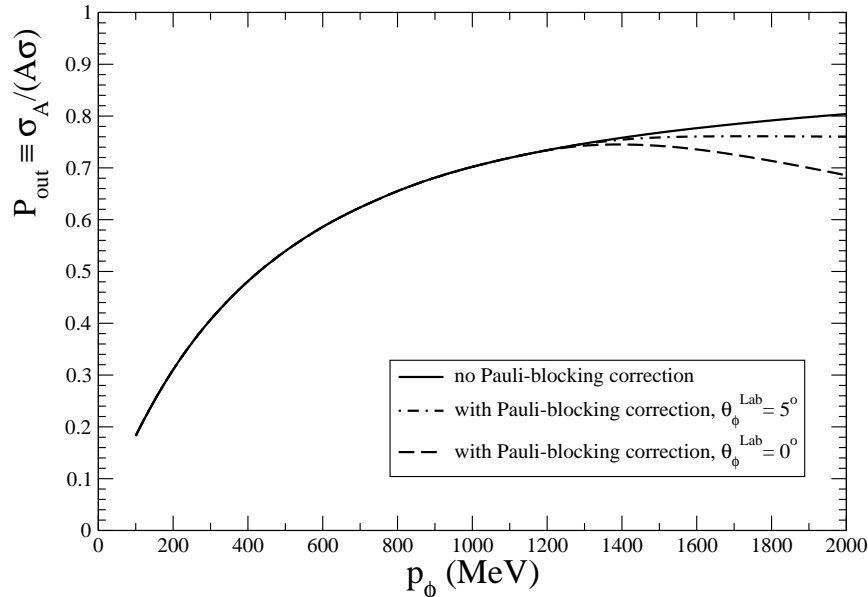


Figure 7: Effect of the Pauli-blocking correction in the $\sigma_A/(A\sigma)$ as a function of the momentum of the ϕ for ^{64}Cu calculated with the CV model.

estimation, using Eq. (15). Note that the Q dependence of this G factor introduces another kinematical variable in the P_{out} observable. We have chosen this extra variable to be the zenithal angle of the ϕ meson in the Lab frame (θ_ϕ^{Lab}). The dashed and dot-dashed curves represent the results considering the Pauli effect for $\theta_\phi^{Lab} = 0^\circ$ and 5° respectively. We can see that significant effects are obtained for large ϕ momenta, what is expectable since large ϕ momenta imply small final nucleon momenta, which are strongly Pauli blocked. The Pauli effect also decreases with the ϕ meson angle because the momentum transfer strongly increases with this angle, therefore the largest effect is obtained for the ϕ forward direction. Since the ϕ photoproduction is strongly forward peaked, one can expect P_{out} to be actually somewhere between the dashed and dot-dashed lines in a real experiment.

In Fig. 8 we show the A dependence of P_{out} , with the CV model, without including the Pauli correction (solid line) and including it for $\theta_\phi^{Lab} = 0^\circ$ and 5° (dashed and dot-dashed lines respectively), all for $p_\phi = 2000$ MeV where the effect is maximum. We observe that the loss of flux due to the Pauli correction is nearly constant in a wide range of the mass number, resulting in a reduction of around 0.1 in P_{out} for the forward direction. Nonetheless this effect is relatively smaller for heavier nuclei compared to the ϕ absorption.

The cross sections calculated in this work are inclusive, summing over all possible final nuclear excited states. The coherent cross section, where the final nucleus is the original one, is not included in the calculation. Its evaluation requires a complete knowledge of the spin and isospin dependence of the elementary $\gamma N \rightarrow \phi N$ amplitude and a different treatment than the one done here. The coherent cross section involves the square of the product of the mass number and the nuclear form factor times the square of the spin-isospin averaged amplitude. The coherent and incoherent cross sections can in principle be

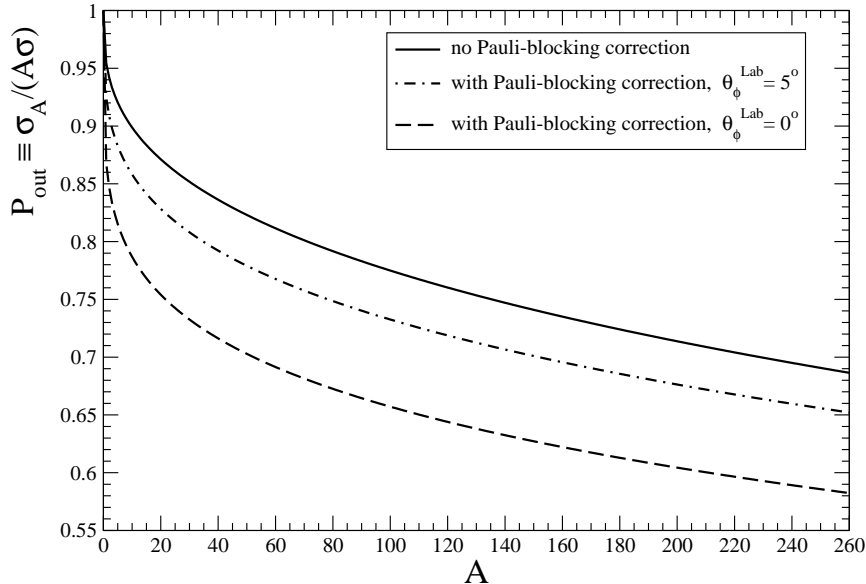


Figure 8: Effect of the Pauli-blocking correction in the $\sigma_A/(A\sigma)$ as a function of the nuclear mass number for $p_\phi = 2000$ MeV calculated with the CV model.

separated experimentally, as done in the case of pions around the delta energy region [33]. However, this might be difficult at higher energies where one usually has a poorer energy resolution. An alternative to this separation, concerning the present work, is to look for ϕ production at finite angles where the momentum transfer to the nucleus is large and the coherent process can be neglected. The incoherent cross section sums the square of the spin-flip and non spin-flip amplitudes, while in the coherent part only the spin independent part contributes. Thus, the ratio $|A F(Q)|^2/(A G(Q))$, where $F(Q)$ is the nuclear form factor and $G(Q)$ the Pauli blocking factor of Eq. (15), is an upper bound for the ratio of coherent to incoherent nuclear cross sections. We have studied this ratio as a function of the photon energy, the ϕ meson angle and the mass number. Our results can be summarized as follows: 1) We find that below $E_\gamma = 2000$ MeV the ratio of coherent to incoherent cross section is smaller than ten percent for all angles and nuclei beyond $A \sim 12$. 2) The ratio increases rapidly with the photon energy, since the momentum transfer decreases. For instance at $E_\gamma = 2400$ MeV, for ^{12}C and forward angles, the ratio can be of the order of 50 percent, but for heavier nuclei the ratio is smaller. For instance, for ^{27}Al this ratio is smaller than 10 percent. We also obtain that beyond 4 degrees the ratio is only a few percent and can be safely neglected. In practical terms, from the experimental point of view, we can say that in nuclei around ^{16}O or heavier and photon energies smaller than 2000 MeV, the coherent contribution is negligible. This discussion may serve to select an experimental setup in which the coherent contribution can be neglected, hence facilitating the interpretation of the data.

5 Conclusions

We have shown that using present models for the ϕ selfenergy in a nuclear medium, conveniently extrapolated to finite ϕ momenta, it is possible to evaluate the survival rate of ϕ produced in nuclear ϕ photoproduction, and how this survival rate is tied to the ϕ width in the medium and the momentum of the ϕ . The survival rates for ϕ coming from photons in the range of 1.6 to 2.4 GeV are of the order of 0.7, a significant deviation from unity, which are measurable experimentally. We have shown the A dependence of the expected $\sigma_A/(A\sigma)$ ratio as well as its dependence on the ϕ momentum. Comparison of the results with experimental numbers of the incoming experiments would help determine the accuracy of the models used. These models could then be used to extrapolate results at other ϕ momenta and one could then get a fair idea of how the ϕ properties are modified in nuclear matter for ϕ at any finite momenta below those studied in the present work.

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