Ohmic Losses in Valence-band Photoemission Experiments

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Abstract

Photoemission experiments involve the motion of an electron near a conducting surface. This necessarily generates heat by ohmic losses from eddy currents. This inelastic scattering of the electrons will result in a downward shift in observed spectra. This effect is most pronounced in poorly conducting metals: in good metals the electron’s field is screened out of the material, while insulators are by definition unable to absorb electromagnetic energy at low frequencies. We give a classification of photoemission processes which shows that the effect is an extrinsic process distinct from final state effects. The shift is illustrated by a model system with a Drude-like conductivity function and a temperature-dependent relaxation time. We give a brief experimental survey of systems in which the ohmic losses may be significant.
I. INTRODUCTION

Photoemission provides the most direct way to measure the energy spectrum of electrons in a solid. The equation given in elementary modern physics texts in the section on the photoelectric effect is

\[ E = E_\nu + \epsilon - W, \]  

where \( E \) is the energy of the emitted electron, \( E_\nu \) is the energy of the incident photon, \( \epsilon \) is the energy of the electron in the solid relative to the chemical potential and \( W \) is the work function. For monochromatic incident light, the spectrum of observed energies \( E \) is the same as that for the electron energies \( \epsilon \) in the solid, except for a known shift.

This equation is known to be oversimplified in a number of respects. For example, an electron can scatter from other electrons on leaving the solid, giving up some of its energy, and even possibly producing secondary electrons which can escape from the sample. In fact, the mean free path of an electron moving with kinetic energies of 20 eV is short, of order 10\( \text{Å} \), while the penetration depth of the light is much longer, meaning that only a small fraction of the electromagnetic energy absorbed actually leads to a photoemitted electron satisfying Eq. 1. Nonetheless, Eq. 1 is used for the analysis of the vast majority of angle-integrated photoemission spectroscopy (AIPES) experiments. The justification for this is the belief that the electrons that absorb a photon fall into two categories: (1) electrons that scatter strongly inelastically, losing an energy of several electron volts, and (2) electrons that do not scatter at all. Electrons in the first category either do not escape from the solid, or they do so with such a low kinetic energy that they can be discarded as uninteresting. Those in the second category are believed to follow Eq. 1. In view of the short mean free path, they must come from a thin layer near the surface, leading to the characterization of photoemission (PE) as a surface-sensitive probe.

Once stated in this way, a possible difficulty with this reasoning suggests itself. The argument depends on having a clean separation of energy scales. If an electron loses an
intermediate amount of energy, the logic breaks down. In modern experiments, spectral features on a scale of 10 - 100 meV are very often of great interest. If many electrons have a substantial probability of undergoing inelastic scattering with that much energy loss, the use of Eq. 1, which forms the basis of the interpretation of nearly all PE data, is no longer correct.

This paper will argue that this separation cannot always be be made, and that the assumptions underlying Eq. 1 can be dangerous in materials with relatively low conductivity or low effective carrier density. This includes many systems of current experimental interest.

Consider the photoemitted electron after it has left the solid at a speed \( v \). Its electric field creates eddy currents and ohmic losses in the solid it is leaving behind. These losses are determined by the dielectric function \( \varepsilon(\omega) \). If the spectral features are of width \( \bar{\hbar}\omega \approx 10 - 100 \text{ meV} \), then it is the Drude (or at least the low-frequency) dielectric response that determines the extent of the low-energy inelastic scattering. Dimensional considerations to be detailed below show that the probability for low energy scattering is

\[
P_1 = \frac{e^2}{v\hbar}A.
\]

Here \( e^2/\hbar v = (1/137)(c/v) \approx 0.8 \) for \( E_\nu = mv^2/2 = 20 \text{ eV} \). \( A \) is a dimensionless number that measures the absorption strength of the material in the Drude regime, but it is generally of order unity for isotropic materials with a Drude-like response. Thus the fraction of electrons that scatter is not small. The important issue is the typical energy loss of the electron. We will estimate below that a typical loss is of order \( \hbar\sqrt{\sigma_0/\tau} \) where \( \sigma_0 \) is the DC conductivity and \( 1/\tau \) is a relaxation rate - the width of the Drude portion of the function \( \sigma(\omega) \). This is in the important 10-100 meV regime for materials in which the effective carrier density is low. This condition is fulfilled in many materials of current interest; in particular a material which undergoes a metal-insulator transition will usually pass through this regime.

Loss processes of the kind treated here, but generally at higher energy, have been considered before. For example, plasmon emission in X-ray PE causes satellite peaks in the measurement of core-level energies [1]. Inelastic losses in EELS (electron energy loss spec-
troscopy) are due to essentially the same physical process [4]. The theoretical description of these experiments is formally similar to that given here. The difference lies in the application of the concepts. The Drude part of the dielectric function is often temperature-dependent. This dependence can be dramatic in certain cases, as, for example, at a superconducting transition or a metal-insulator transition. This means that the extrinsic losses are also temperature-dependent. It is then not permissible to interpret temperature-dependent changes in an AIPES spectrum as being due only to changes in the density of states (DOS). The Drude part may also be sample-dependent. Thus, changes in AIPES from sample to sample may reflect changes in the conductivity, as well as changes in the DOS. Such effects are not so likely to occur in the plasmon part of the dielectric function. The low-frequency end of these theories has, however, never been of great interest before for PE. It is only the great advances in PE resolution, now \( \leq 20 \) meV for many experiments, which lends new interest to the application of these known concepts. A short summary of the results has appeared in previous publications [3].

In this paper we first give a classification of PE processes in order to situate the process treated here in the context of the overall theory of PE. Next, a detailed derivation of the loss function is presented, and its effect on observed spectra is illustrated in a simple case. Finally we give a short survey of some experimental systems that are candidates for the effect.

II. CLASSIFICATION OF LOSS PROCESSES

The appropriate formalism is that due to Schaich and Ashcroft [4] and Caroli et al. [5]. We present here a graphical interpretation of this formalism.

The conductivity (absorption of a low energy photon) in a solid is described by a current-current correlation. In the language of Feynman diagrams, we have a line representing an electron and a line representing a hole. Together, these make a loop with two vertices, the dipole matrix elements, at the ends. The PE experiment also involves the absorption of a
photon, but there are two important differences between a conductivity diagram and a PE diagram for angle-resolved photoemission (ARPES). First, there is a measurement of the energy and momentum of the electron. Second, the electron line corresponds to a vacuum state, not a bulk state.

We represent the measurement by a rectangle inserted in the electron line. (In order to give formal expressions corresponding to each diagram, it is necessary to use the Keldysh rather than the Feynman formulation of perturbation theory, but the distinction does not affect the qualitative considerations given here.) This may now be called a current-current-current correlation function, as the electron line has been split into two. (The nomenclature is somewhat misleading since it would seem to imply that PE measures a third-rank tensor. It does not.)

The simplest diagram for a PE process is shown in Fig. 1. The electron propagates freely to the detector in this case. The measurement fixes the energy and momentum of the electron line, so it contributes only a delta function to the diagram and what we have is the textbook process of Eq. 1. If the momentum is averaged over, then the experiment is AIPES, which is the experiment we will concentrate on in this paper.

Fig. 2 is the most-studied of all the diagrams. The thick solid line is the full propagator for the hole. This diagram represents intrinsic many-body effects. All many-body effects on the propagation of a hole in the bulk system must be included, in principle. The reason for calling this diagram “intrinsic”, is that nothing about the photoemission process itself is involved. The emitted electron still propagates freely to the detector. The diagram of Fig. 2 includes that of Fig. 1. To the extent that only these diagrams are important, PE measures the spectral function of a hole. In simpler language, it measures the occupied DOS.

Fig. 3 is the effect of the interaction of the emitted electron with the hole left behind. This is usually called a final state interaction, an ”extrinsic” process, as it depends on the PE process itself. The rate of such interactions decreases with incoming photon energy, and this diagram is usually neglected. This is called the sudden approximation. Fig. 4 represents an extrinsic process in the Born approximation. The thick interaction line is fully screened,
and thus includes interactions with all the electrons in the material. This changes the energy and the momentum of the observed particle. This diagram is the subject of the present work. Note that the distinction between the final state interactions and the extrinsic inelastic loss processes is a clean one. The two should not be confused. The ohmic losses are, roughly speaking, due to interactions with the dynamic image charge, not the photohole.

This is not a complete listing of all possible processes, as one can also have scattering from impurities of the photoelectron, reflection at the surface, and so on. We have tried to give a reasonable overview of the possibilities that make an interesting comparison or contrast with the ohmic process.

III. DERIVATION FOR CUBIC SYSTEMS

We begin with the general problem of a sample that occupies the half-space \( z < 0 \). There is a time-dependent charge density \( \rho(\vec{r}, t) \) outside it. The Fourier component at frequency \( \omega, \rho(\vec{r}, \omega) \), sets up an electric field for \( z < 0 \) given by

\[
\vec{E}(\vec{r}, \omega) = \int_{z' > 0} \rho(\vec{r}', \omega) \frac{\vec{r}' - \vec{r}}{|\vec{r}' - \vec{r}|} \frac{d^3 r'}{1 + \epsilon(\omega)}. \tag{3}
\]

This result follows from the usual image-charge calculation, namely the solution of Maxwell’s equations with the standard surface boundary conditions. Nevertheless, there are certain approximations. The normal skin effect is assumed because the wavevector dependence of \( \epsilon \) is neglected. This could break down for clean materials or at quite high frequencies. Fortunately, both of these conditions are outside of the applications which interest us here. Secondly, the real surface is not of infinitesimal thickness, and corrections for this have been neglected. Finally, we have specialized to the case that the material has cubic symmetry: the electric field and the electric displacement at any given frequency are simply proportional. The field drives currents that generate heat. The total energy thus produced is:

\[
Q = \frac{1}{2} \int_{z < 0} d\omega \Re \sigma(\omega) |\vec{E}(\vec{r}, \omega)|^2 d^3 r. \tag{4}
\]
These results are quite general for any charge density for which the integral in Eq. 4 is convergent. Now we specialize to the PE process, which is modeled as a charge generated at the surface at time \( t = 0 \) leaving the solid at an angle \( \theta \) to the normal. The charge density is

\[
\rho(\vec{r}, t) = -e\Theta(t)\delta(z - vt \cos \theta)\delta(y)\delta(x - vt \sin \theta),
\]

(5)

where \( \Theta(t) \) is the step function. Since in most experiments \( \theta \leq \pi/10 \), we set \( \cos \theta = 1 \) henceforth. The Fourier transform is

\[
\rho(\vec{r}, \omega) = \frac{-e}{2\pi v \cos \theta}\delta(x)\delta(y)e^{-i\omega z/v}\Theta(z).
\]

(6)

Using Eq. 3, we have

\[
\vec{E}(\vec{r}, \omega) = \frac{-e}{2\pi v \cos \theta} \frac{2}{1 + \epsilon(\omega)} \int_0^\infty dz' e^{-i\omega z'/v \cos \theta} \frac{\vec{r} - z'\hat{z}}{|\vec{r} - z'\hat{z}|^3}
\]

(7)

as the field inside the sample. Substituting this into Eq. 4 yields

\[
Q = \frac{2e^2}{\pi v} C \int_0^\infty d\omega \frac{\text{Re} \sigma(\omega)}{\omega|1 + \epsilon(\omega)|^2},
\]

(8)

where \( C \) is a dimensionless integral:

\[
C = \int_0^\infty \rho d\rho \int_0^\infty dz \left( \rho^2 |I_1(\rho, z)|^2 + |I_2(\rho, z)|^2 \right),
\]

(9)

with

\[
I_1 = \int_0^\infty \frac{e^{-i\zeta} d\zeta}{[\rho^2 + (z + \zeta)^2]^{3/2}}
\]

(10)

and

\[
I_2 = \int_0^\infty \frac{e^{-i\zeta}(z + \zeta) d\zeta}{[\rho^2 + (z + \zeta)^2]^{3/2}}.
\]

(11)

\( C \) may be estimated as about 2.6.

This classical calculation yields not only the total energy loss but also the loss into each frequency interval \( d\omega \). Quantum-mechanically, this is interpreted as the inelastic energy.
loss at this energy due to a scattering event from the bulk electrons. It was calculated assuming no change in the trajectory of the electron, and is therefore essentially the Born approximation. Fig. 4 could also have several scattering lines. This would correspond to approximations beyond Born. When such processes are important, then the amount of information contained in a PE spectrum decreases, just as in the case of multiphonon events in neutron scattering (for example).

The relative differential probability to lose energy $\hbar \omega$ is

$$P(\omega) = \frac{2e^2C}{\pi \hbar \nu \omega^2} \frac{\mathbb{R} \sigma(\omega)}{|1 + \epsilon(\omega)|^2}. \quad (12)$$

There is also a probability for forward scattering $P_0$. This is the probability that an electron loses zero energy. Including this possibility, the observed intensity is given by

$$I(\omega, T) = P_0(T)N(\omega)f(\omega) + \int_{-\infty}^{\infty} P(\omega - \omega', T)N(\omega')f(\omega')d\omega', \quad (13)$$

Here $N(\omega)$ is the DOS and $f$ is the Fermi function. The total normalization is given by our assumption that there is either one or zero scatterings:

$$1 = P_0 + \int_{0}^{\infty} P(\omega)d\omega. \quad (14)$$

Equations having the form of Eq. 13 are well-known in considerations of background in PE. For example, the Shirley background would result from the phenomenological assumption that $P(\omega)$ is a constant to be fit to experiment.

**IV. EXPERIMENTAL CONSEQUENCES**

From the point of view of experiment, one would like rules of thumb for recognizing the situations in which the ohmic effect is likely to be important. To this end, we investigate more closely the function $P(\omega)$. Note first that

$$P(\omega) = \frac{2e^2C}{\pi \hbar \nu \omega^2} \frac{\mathbb{R} \sigma(\omega)}{|1 + \epsilon(\omega)|^2}. \quad (15)$$
vanishes at small frequencies if the system is a true insulator, since $\Re \sigma$ vanishes faster than any power of $\omega$. This must be true from a physical standpoint: an insulator cannot absorb electromagnetic energy. If the system is a very good metal, then $\sigma_0 \equiv \Re \sigma(\omega = 0)$ is large, and at low frequencies $\epsilon(\omega) \to 4\pi i \sigma_0 / \omega$ leading to

$$P(\omega) \to \frac{1}{16\pi^2 \sigma_0} \frac{2e^2 C}{\pi \hbar v}.$$  \hspace{1cm} (16)

The vanishing of $P(\omega)$ in this limit is also required by elementary physics. A perfect metal must completely screen out the field at sufficiently low frequencies. We conclude that materials in the intermediate range are of most interest. Certainly if PE is done across a metal-insulator transition this range is explored.

There is no general formula that determines the importance of ohmic losses, because $\epsilon(\omega) = 1 + 4\pi i \sigma / \omega$ can take many forms, with the consequence that even materials with similar DC conductivities may give quite different results. Furthermore, the effect depends on the frequency range which is being explored. However, it very often happens that in the low-frequency regime the two-parameter Drude function

$$\sigma(\omega) = \frac{\sigma_0}{1 - i\omega \tau}.$$ \hspace{1cm} (17)

is a reasonable approximation. In this paper, we will focus on this particular form. Even for materials for which the Drude form is not a very good quantitative approximation for the conductivity, it may still be true that one parameter for the overall scale and one for the width are sufficient. In this case, the results below should give a reasonable guide.

The Drude form for $P(\omega)$, obtained by substituting Eq. 17 into Eq. 12 and rearranging, is:

$$P(\omega) = \frac{2e^2 C \sigma_0 \tau^2}{\pi \hbar v} \left\{ \left[ \omega^2 \tau^2 - (2\pi \sigma \tau - 1/2) \right]^2 + 2\pi \sigma_0 \tau - 1/4 \right\}^{-1}.$$ \hspace{1cm} (18)

The first thing to note about this function is its normalization. Integrating Eq. 18 we find

$$\int_0^\infty P(\omega) d\omega = \frac{e^2 C}{4\pi \hbar v}.$$ \hspace{1cm} (19)
In terms of the notation of Eq. 4, this says \( A = C/4 \approx 0.65 \). Substituting in a velocity corresponding to the energy \( E_\nu = 20 \text{ eV} \) we find \( P_1 = 0.5 \). Hence the integrated effect is substantial. Since this result is independent of the parameters, it is a generic result not limited to the particular details of the Drude form. Bear in mind that this is not the only contribution to \( P_1 \) since losses at higher energies, for example losses due to interband processes, will also contribute. On the other hand, the fact that \( P_1 \) is comparable to unity at typical experimental energies also implies that we near the limit of validity of the Born approximation.

Written in terms of the variable \( \omega^2 \tau^2 \), \( P(\omega) \) is a Lorentzian centered at \( 2\pi \sigma \tau - 1/2 \) and width \( \sqrt{2\pi \sigma_0 \tau - 1/4} \). In terms of the physical variable \( \omega \), we have two regimes. If \( 2\pi \sigma_0 \tau < 1/2 \), then the function is peaked at \( \omega = 0 \), and falls off at large frequencies as \( \omega^{-4} \) on a characteristic scale of \( 1/\tau \). If \( 2\pi \sigma_0 \tau > 1/2 \), then the function is peaked at

\[
\omega_{\text{peak}} = \sqrt{2\pi \sigma_0 / \tau - 1/2} \approx \sqrt{2\pi \sigma_0 / \tau},
\]

and the width of the peak is

\[
\Delta \omega = \frac{1}{\tau} (2\pi \sigma_0 \tau - 1/4)^{1/4} \rightarrow \left( \frac{2\pi \sigma_0 \tau}{\tau^3} \right)^{1/4}.
\]

The limiting behaviors for \( 2\pi \sigma_0 \tau \gg 1 \) are also given.

The overall behavior of \( P(\omega) \) is illustrated in Fig. 5. It is plotted for a fixed value of \( \sigma_0 / \tau \). In the Drude model

\[
\frac{\sigma_0}{\tau} = \frac{ne^2}{m^*} = \frac{\omega_p^2}{4\pi},
\]

where \( \omega_p \) is an effective plasma frequency. Thus, in a system where the effective carrier concentration is fixed and the temperature dependence comes either from the relaxation time or from \( \sigma_0 \), Fig. 5 represents the evolution of \( P(\omega) \) with temperature. This graph is for a low effective carrier density: \( \sigma_0 / \tau = 10^4 \text{ meV}^2 \), which corresponds to an effective plasma frequency of 354 meV. The relaxation times are in units of meV\(^{-1}\) where \( \hbar = 1 \), that is \( 6.6 \times 10^{-13} \text{s} \). As \( \tau \) increases, the peak position moves from \( \omega = 0 \) to the limiting value \( \sqrt{2\pi \sigma_0 / \tau} \), while it continually sharpens up.
These structures in $P(\omega)$ manifest themselves in the observed intensity according to Eq. 13. Consider the case of zero temperature. Then $f(\omega)$ becomes $\theta(-\omega)$ and we have:

$$I(\omega) = N(\omega)\theta(-\omega)P_0 + \int_0^\infty P(\omega' - \omega)N(\omega')d\omega'.$$

We are only interested in the second term, the inelastic component, here. Furthermore, we shall take a constant DOS $N(\omega) = N_0 = 1$ so that the features in the observed intensity come only from the energy dependence of the scattering. Thus we normalize $P(\omega)$ to unity.

$$I_{\text{inelastic}}(\omega) = \int_0^\infty P(\omega' - \omega)d\omega'$$

The integral rearranges the apparent DOS. The DOS will appear unchanged at frequency $\omega$ if $|\omega| > \omega_{\text{peak}}$ where $\omega_{\text{peak}} = 0$ or $\sqrt{2\pi\sigma_0\tau} - \frac{1}{2\pi^2 \tau}$ depending on the size of $\sigma_0\tau$. The photoelectrons observed at large $|\omega|$ are shifted, but it appears as if the (flat) DOS is unaffected at these frequencies. At low frequencies the situation is more complicated. The basic scenario for a constant $\sigma_0/\tau$ (effective carrier density) is this: at 'small' $\tau << \tau_c$ $P(\omega)$ is peaked around $\omega = 0$ so the PE loss is negligible. At the crossover lifetime $\tau_c = \sqrt{\frac{m^*}{4\pi ne^2}}$ the peak of $P(\omega)$ begins to shift to a higher frequency. In addition, $P(\omega)$ is very spread out for $\tau$ in this region, so the apparent loss becomes large and may appear pseudogap. When $\tau >> \tau_c$ the shape of the observed intensity $I(\omega)$ looks unchanged, but the entire curve is shifted downwards by an amount comparable to the effective plasma frequency, and this may look like a gap. This evolution of $I(\omega)$ with increasing $\tau$ is shown in Fig. 6. Bear in mind that this is only the inelastic component and it must be combined with the elastic intrinsic part. This may involve fitting since the overall normalization of the elastic part ($P_0$) would normally be hard to determine.

It should be noted that while most of the peak shift of $P(\omega)$ occurs over a relatively small range of $\tau$ the evolution of $I(\omega)$ takes place over several orders of magnitude of $\tau$ since $P(\omega)$ is spread out over a large frequency range. We can see this by plotting the shift of the intensity curve vs. $\tau$ as in Fig. 4. The shift is defined by the position in energy of the point where the intensity is half its maximum, a popular definition of the gap in PE experiments.
V. SURVEY OF EXPERIMENTAL SYSTEMS

Experimentally, the best candidates for gap-like structures are the one-dimensional conductors. There are several one-dimensional materials that show the development of some sort of a gap in the spectrum without corresponding observations in other properties. We mention here only two.

TTF-TCNQ is a one-dimensional organic chain compound. It has a $2k_F$ charge-density-wave transition at $T_F = 54K$ with an energy gap $2\Delta \sim 40 \text{ meV}$, as measured by the activation energy of transport and thermodynamic quantities. The photoemission spectrum is very much at odds with the transport and thermodynamics. At 70K, there is a gap of about $0.15 \sim 0.2eV$, judging by the peak in an energy scan at an angle that corresponds to $k_F$. This peak does not move toward the Fermi energy as T increases. Rather, the spectrum just broadens out, but without increasing the weight at the Fermi energy. In fact, the weight if anything decreases. This material is a good candidate for a pseudogap of partly extrinsic origin.

$(\text{TaSe}_4)_2\text{I}$ presents a somewhat similar scenario. It has a $2k_F$ charge density wave transition at about 250K. A number of its properties: resistivity, magnetic susceptibility, and optical conductivity can be explained by assuming that the gap is roughly 250 meV, and remains at about this value, independent of temperature, even after long-range positional order is lost. But the gap in photoemission is about 500 meV, independent of temperature. Here we may have a gap which is split about 50-50 between intrinsic and extrinsic.

The colossal magnetoresistance manganate materials are good candidates for extrinsic losses in the neighborhood of their metal-insulator transition. The cubic material
La_{0.67}Ca_{0.33}MnO_3 is in fact the only system for which model calculations of the extrinsic losses have been carried out [3]. The material has a metal-insulator transition at 260 K. In the 'metallic' state at 80 K, there is a strong negative slope in \( I(\omega) \) for at least 0.6 eV below \( \mu \) [12]. There is a sharp break in slope at \( \mu \), presumably indicative of a nonzero density of states at \( \mu \). But in the insulating state at 280 K there appears to be no Fermi edge at all - the observed intensity is flat at \( \mu \) and weight has moved back from \( \mu \). These features might be taken as indicating the presence of a pseudogap which opens in the insulating state. However, this movement of spectral weight can be produced by extrinsic effects, as shown in [3].

The layered material La_{1.2}Sr_{1.8}Mn_2O_7, on the other hand, has been shown by Dessau and Saitoh [13] to have a momentum-dependent pseudogap somewhat similar to the high-T_c materials, but rather larger, in the range of 200 meV or so at the maximum. Again, the momentum dependence indicates that the pseudogap is mostly of intrinsic origin. The layered manganate materials are an interesting case because of their extremely high resistivities. La_{1.2}Sr_{1.8}Mn_2O_7 itself has an a-b plane resistivity of about 6 m \( \Omega \)-cm at \( T = 50 \) K, which is in the 'metallic' phase. This is an extraordinarily high value. It is so close to being insulating that one would probably not expect ohmic effects [14].

Another class of poorly conducting materials are the Kondo insulators. This class of systems shows a resistivity dominated by the Kondo effect at higher temperatures crossing over to activated behavior at low temperatures, with absolute values which can be quite high. The resistivity as a function of temperature is monotonic, in contrast to the CMR systems. For example, the system YbB_{12} has such a crossover at about 50K. This material also shows some indications of extrinsic processes. Even on the metallic side, the observed density of states slopes downward strongly to the Fermi energy [13]. As the resistivity increases, there is a further depression of weight near \( E_F \), with finally a gap-like structure of about 20 meV appearing at 6 K.

The best-known pseudogap of all is that in the high-temperature superconductor Bi_{2}Sr_{2}CaCu_{2}O_{8} [11]. In this case, the pseudogap is both momentum- and temperature-
dependent, appearing mainly along the Cu-O bond direction. This is a strong indication that the pseudogap is largely intrinsic, as the extrinsic losses should not have a strong momentum dependence. Bi$_2$Sr$_2$CaCu$_2$O$_8$ is roughly tetragonal and the theory developed for the cubic case is not applicable. Its conductivity in the a-b plane is fairly high. The c-axis conductivity, which is very low, also plays a role, however. The theory for such a strongly anisotropic system has not been worked out in detail, and the combined effect of $\epsilon_{xx}$ and $\epsilon_{zz}$ is not yet completely clear. It seems most likely that lineshapes are affected by extrinsic effects, but that the size of the pseudogap is correctly given by the naive analysis.

The high-temperature superconductors and related cuprates are a class of systems for which some energy-dependent PE has been done. This is very interesting from the point of ohmic effects, which are predicted to decline as $1/v \sim 1/\sqrt{E_\nu}$, as seen from Eq. 12. For example, Sr$_2$CuO$_2$Cl$_2$ has been investigated using ARPES for $E_\nu = 22$eV and $E_\nu = 35$eV [16] and Pb-doped Bi$_2$Sr$_2$CaCu$_2$O$_8$ at $E_\nu = 32$eV, $E_\nu = 40$eV and $E_\nu = 50$eV [17]. In every case, changes in intensity as a function of $E_\nu$ are seen. In the case of Pb-doped Bi$_2$Sr$_2$CaCu$_2$O$_8$ they go in the direction of increasing binding energy as $E_\nu$ is decreased, which is in qualitative agreement with the present scenario. In addition, the lineshapes are rather broad and also energy dependent.

It is evident from this brief survey that an important direction for development of the theory would be a generalization to tetragonal and orthorhombic materials, as relatively few of the most interesting systems are cubic. The image calculation can be performed analytically for the tetragonal case if the emission is along the c-axis, but the loss function is rather complicated and will be published elsewhere. The orthorhombic case is even more complicated, as the image method is not applicable.

VI. CONCLUSION

Changes in the density of states of electronic systems on the scale of 10 to 100 meV are of great interest in today’s PE experiments. These energies often coincide with the scale of
relaxation rates and conductivities or combinations thereof. This means that ohmic losses may affect the interpretation of the data. This effect is distinct from others, such as final state effects, that have been considered in detail in the past.

General formulas for the loss in terms of the dielectric function may be given. In the case where a Drude form is appropriate, we may characterize the situation as follows. The inelastic loss will be most pronounced if the lifetime is roughly \( \tau_c \approx 1/(2\pi\sigma_0) \). The magnitude of the loss is on the scale of the maximum peak shift \( \omega_{\text{peak}} = \sqrt{2\pi\sigma_0/\tau} \) near the effective plasma frequency. If the effective plasma frequency is large, as in a good metal, then the inelastically scattered electrons are emitted well below the chemical potential as secondaries, and the elastic peak is unsullied. If the product \( \sigma_0\tau \) is very small, as in a near-insulator, the scattered electrons are emitted with very small loss and cannot be distinguished from the unscattered ones. In the intermediate case, however, a combination of the elastic and inelastic terms is required to give the observed spectrum. This ohmic effect may account for the observation of a gap or pseudogap in some systems with relatively low intrinsic conductivities in one or more directions. In general, however, low conductivity is a necessary, not a sufficient, condition, for the ohmic effects to be important. The effective carrier density and the frequency range of the measurements must also be taken into account.

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FIGURES

FIG. 1. The simplest diagram for the calculation of photoemission intensities. A photon (wavy line) of energy $E_\nu$ is absorbed by the system, creating an electron in a vacuum state (right-going line) with energy $E$ and a non-interacting hole in a bulk state (left-going line) with energy $\epsilon$. The measurement of the electron energy and momentum is represented by the dark rectangle, and distinguishes this diagram from an ordinary conductivity diagram. If this diagram is the only important one, the measured intensity in AIPES can be interpreted as the non-interacting DOS.

FIG. 2. Modification of the simplest diagram to include bulk interactions of the hole, indicated by the thick line. This diagram includes the previous one. If this diagram is the only important one, the measured intensity in AIPES can be interpreted as the interacting DOS.

FIG. 3. The diagram for final state effects in which the hole interacts with the outgoing electron.

FIG. 4. The diagram for scattering of the outgoing electron from other electrons in the material. This includes the extrinsic ohmic process.

FIG. 5. Evolution of the relative differential probability $P(\omega)$ for losing an energy of $\hbar\omega$. For the different plots, the effective carrier density is held fixed: $\sigma_0/\tau = 10^4$ meV$^2$, while the relaxation times are in units of meV$^{-1}$ where $\hbar = 1$, that is $6.6 \times 10^{-13}$ s.

FIG. 6. The observed inelastic intensity corresponding to the loss functions in Fig. 5.

FIG. 7. The shift of the inelastic intensity curve vs. $\tau$ as in Fig. 5. The shift is defined by the position in energy of the point where the intensity is half its maximum.
