

# **The DMFT object of the CP-PAW code**

Peter E. Blöchl

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Institute of Theoretical Physics; Clausthal University of Technology;  
D-38678 Clausthal Zellerfeld; Germany;  
<http://www.pt.tu-clausthal.de/atp/>

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# Chapter 1

## Purpose and theoretical background

The purpose of the DMFT object is to prepare an interface to the solver for a quantum impurity in the context of dynamical mean-field theory. The derivations follow closely the paper of Blöchl, Pruschke and Potthoff[1], in the following referred to **BPP**.

### 1.1 Grand potential and density-matrix functional as starting point

The **grand potential** has the form[1]

$$\Omega_{\beta,\mu}^{KB}[\hat{h} + \hat{W}] = \min_{|\psi_n\rangle, f_n \in [0,1]} \text{stat}_{\Lambda} \left\{ \sum_n f_n \langle \psi_n | \hat{h} | \psi_n \rangle + \tilde{F}_{\beta}^{\hat{W}} \left[ \sum_n |\psi_n\rangle f_n \langle \psi_n| \right] - \mu \sum_n f_n - \sum_{m,n} \Lambda_{m,n} \left( \langle \psi_n | \psi_m \rangle - \delta_{m,n} \right) \right\} \quad (1.1)$$

where the **density-matrix functional**  $\tilde{F}_{\beta}^{\hat{W}}$  is expressed[1] by the **Luttinger-Ward functional** [2]  $\Phi^{LW}[\mathbf{G}, \hat{W}]$  as

$$\begin{aligned} \tilde{F}_{\beta}^{\hat{W}}[\rho] = & \frac{1}{\beta} \text{Tr} \left[ \rho \ln(\rho) + (\mathbf{1} - \rho) \ln(\mathbf{1} - \rho) \right] \\ & + \text{stat}_{h'} \text{stat}_{\mathbf{G}, \Sigma} \left\{ \Phi_{\beta}^{LW}[\mathbf{G}, \hat{W}] - \frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \ln \left[ \mathbf{1} - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \right\} \right. \\ & \left. + \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \mathbf{G}(i\omega_{\nu}) - \left[ \mathbf{G}(i\omega_{\nu}) - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left( \mathbf{h}' - \bar{\mathbf{h}} \right) \right\} \end{aligned} \quad (1.2)$$

Here,  $|\psi_n\rangle$  are one-particle wave functions. They play the role of **natural orbitals**, the eigenstates of the **one-particle density matrix**  $\rho$ .  $f_n$  are the occupations, the eigenvalues of the one-particle density matrix. The orthonormality of the natural orbitals is enforced with the method of Lagrange multipliers. The Lagrange multipliers are  $\Lambda$ . The chemical potential is  $\mu$  and  $\beta = 1/(k_B T)$ . The full Hamiltonian consists of a non-interacting part  $\hat{h} = \frac{\hat{p}^2}{2m_0} + \hat{v}_{ext}$  and an electron-electron interaction  $\hat{W}$ .

With  $\text{stat}_x Y$ , we denote the stationary condition, which requires that the partial derivatives of  $Y$  with respect to the variable  $x$  is obeyed. It explicitly does not require the stationary point to be an extremum.

$G(i\omega)$  is the Green's function and  $\Sigma(i\omega_\nu)$  is the self energy. The Matsubara sum runs over the **Matsubara frequencies** (see appendix A on p. 38)

$$\omega_\nu = (2\nu - 1) \frac{\pi}{\hbar\beta} \quad \text{for } \nu \in \mathbb{Z} \quad (1.3)$$

$\mathbf{h}'$  are the Lagrange multipliers for the **density-matrix constraint**.

$$\rho = \frac{1}{\beta} \sum_{\nu} e^{i\beta\hbar\omega_\nu 0^+} G(i\omega_\nu) \quad (1.4)$$

The Hamiltonian  $\bar{\mathbf{h}}$  is a non-local Hamiltonian directly related to the one-particle density matrix via (see Eq. BPP24)

$$\bar{\mathbf{h}} = \mu \mathbf{1} + k_B T \ln \left[ \frac{\mathbf{1} - \rho}{\rho} \right] \quad (1.5)$$

In practice, we perform calculations at fixed particle number. Therefore, we calculate the **Helmholtz potential** by adding the term with the particle number<sup>1</sup>  $N_{\beta,\mu}$  to the grand potential.

$$\begin{aligned} A_{\beta,N}[\hat{h} + \hat{W}] &= \text{stat}_{\mu} \left\{ \Omega_{\beta,\mu} + \mu N_{\beta,\mu} \right\} \\ &= \min_{|\psi_n\rangle, f_n \in [0,1]} \text{stat}_{\mu, \Lambda} \left\{ \sum_n f_n \langle \psi_n | \hat{h} | \psi_n \rangle + Q_{\beta}^{\hat{W}} \left[ \sum_n |\psi_n\rangle f_n \langle \psi_n| \right] \right. \\ &\quad \left. + \frac{1}{\beta} \sum_n \left[ f_n \ln(f_n) + (1 - f_n) \ln(1 - f_n) \right] \right. \\ &\quad \left. - \mu \left[ \sum_n f_n - N \right] - \sum_{m,n} \Lambda_{m,n} \left( \langle \psi_n | \psi_m \rangle - \delta_{m,n} \right) \right\} \end{aligned} \quad (1.6)$$

where  $Q_{\beta}^{\hat{W}}$  is the density-matrix functional without the entropy contribution of a non-interacting electron gas with the same density matrix. The entropy contribution is taken care of with the **Mermin functional**[3] to describe DFT calculations with electrons at finite temperature.

$$\begin{aligned} Q_{\beta}^{\hat{W}}[\rho] &= \tilde{F}_{\beta}^{\hat{W}}[\rho] - \frac{1}{\beta} \text{Tr} \left[ \rho \ln(\rho) + (\mathbf{1} - \rho) \ln(\mathbf{1} - \rho) \right] \\ &= \text{stat}_{\mathbf{h}'} \text{stat}_{\mathbf{G}, \Sigma} \left\{ \Phi_{\beta}^{LW}[\mathbf{G}, \hat{W}] - \frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \ln \left[ \mathbf{1} - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \right\} \right. \\ &\quad \left. + \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \mathbf{G}(i\omega_{\nu}) - \left[ \mathbf{G}(i\omega_{\nu}) - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left( \mathbf{h}' - \bar{\mathbf{h}} \right) \right\} \end{aligned} \quad (1.7)$$

## 1.2 Separate out static correlations: Towards DFT+

In appendix D it is shown that the correlation contribution  $Q_{\beta}^{\hat{W}}$  can be divided into one part having a frequency-dependent self energy and several contributions having a frequency-independent self

<sup>1</sup>  $N_{\beta,\mu} = \sum_n f_n$  is the particle number of the equilibrium ensemble defined by a temperature and the chemical potential.

energy. The Luttinger-Ward functional with frequency-independent self energy can be expressed by the density matrix alone, that is, the Green's function enters always in the form of the integral

$$\frac{1}{\beta} \sum_{\nu} \mathbf{G}(i\omega_{\nu}) e^{i\hbar\omega_{\nu}\beta 0^+} \quad (1.8)$$

This notion allows us to split the interaction contribution  $Q^{\hat{W}}$  into different parts as

$$Q^{\hat{W}}[\rho] \approx Q^{DFT, \hat{W}}[\rho] - Q^{DFT, \hat{W}_1}[\rho] + Q^{HF, \hat{W}_1}[\rho] - Q^{HF, \hat{W}_2}[\rho] + Q^{\hat{W}_2}[\rho] \quad (1.9)$$

of which only the last term has a frequency-dependent self energy.

In order to keep the computational effort at a reasonable level, we consider three different interactions, namely the full interaction  $\hat{W}$ , a simplified interaction  $\hat{W}_1$  used for static correlations, and interaction  $\hat{W}_2$ , which is even further simplified, and which is used for dynamic correlations. The full interaction is treated on the most simple level of the description, namely DFT, while the most complex many-particle treatment is limited to the most approximate interaction, namely  $\hat{W}_2$ .

With a density-functional contribution  $Q^{DFT, \hat{W}}[\hat{\rho}]$  acting as double-counting term and a screened Hartree-Fock contribution  $Q^{HF, \hat{W}}[\hat{\rho}]$  described later, we obtain

#### HELMHOLTZ POTENTIAL

$$\begin{aligned} A_{\beta, N}[\hat{h} + \hat{W}] = & \min_{|\psi_n\rangle, f_n \in [0,1]} \text{stat}_{\mu, \Lambda} \left\{ \sum_n f_n \langle \psi_n | \frac{\hat{p}^2}{2m_e} | \psi_n \rangle + \int d^3r n(\vec{r}) v_{\text{ext}}(\vec{r}) \right. \\ & + \frac{1}{2} \int d^3r \int d^3r' \frac{e^2 n(\vec{r}) n(\vec{r}')}{4\pi\epsilon_0 |\vec{r} - \vec{r}'|} + E_{xc}[n] + \frac{1}{\beta} \sum_n [f_n \ln(f_n) + (1 - f_n) \ln(1 - f_n)] \\ & + \underbrace{Q_{\beta}^{HF, \hat{W}_1}[\rho] - Q_{\beta}^{DFT, \hat{W}_1}[\rho, n]}_{\text{local hybrid functional}} + \underbrace{Q_{\beta}^{\hat{W}_2}[\rho] - Q_{\beta}^{HF, \hat{W}_2}[\rho]}_{Q_{\beta}^{\text{dyn}, \hat{W}_2}[\rho]} \\ & \left. - \mu \left[ \sum_n f_n - N \right] - \sum_{m,n} \Lambda_{m,n} \left( \langle \psi_n | \psi_m \rangle - \delta_{m,n} \right) \right\} \end{aligned} \quad (1.10)$$

where

$$\begin{aligned} n(\vec{r}, \sigma, \sigma') &= \sum_n \langle \vec{r}, \sigma | \psi_n \rangle f_n \langle \psi | \vec{r}, \sigma' \rangle \\ n(\vec{r}) &= \sum_{\sigma} n(\vec{r}, \sigma, \sigma) \\ \rho_{a,b} &= \sum_n \langle \pi_a | \psi_n \rangle f_n \langle \psi | \pi_b \rangle \end{aligned} \quad (1.11)$$

We distinguish two different corrections:

1. a screened Hartree-Fock correction with an interaction  $\hat{W}_1$ .

The Hartree-Fock term  $Q_{\beta}^{HF,\hat{W}}$  is equal to  $Q_{\beta}^{\hat{W}}$  when only the first-order term of the Luttinger-Ward functional in the interaction is considered. It is obtained as (BPP-Eq.43)

$$Q_{\beta}^{HF,\hat{W}}[\rho] = \frac{1}{2} \sum_{a,b,c,d} U_{a,b,d,c} [\rho_{d,a}\rho_{c,b} - \rho_{c,a}\rho_{d,b}] \quad (1.12)$$

(See also appendix E on p.57.)

The first term,  $E_H$ , is the **Hartree energy** and the second term,  $E_X$ , is the **exchange energy**

$$\begin{aligned} E_H &= \frac{1}{2} \sum_{a,b,c,d} U_{a,b,d,c} \rho_{d,a} \rho_{c,b} \\ E_X &= -\frac{1}{2} \sum_{a,b,c,d} U_{a,b,d,c} \rho_{c,a} \rho_{d,b} \end{aligned} \quad (1.13)$$

For the Hartree-Fock contribution we restrict the U-tensor to on-site contributions and certain bond-terms. In the spirit of the hybrid functionals or the GW method, we scale the U-tensor. The neglect of U-tensor elements that do not reside on an atom pair is motivated by the fact that screening becomes stronger with increasing distance. Part of this approximation is to neglect the interactions not captured by the local orbital basis. This term is evaluated in the `paw_lmto`-object.

2. a **DFT double counting term**  $Q_{\beta}^{DFT,\hat{W}_1}$ :

$$Q_{\beta}^{DFT,\hat{W}}[\hat{\rho}] = \frac{1}{2} \int d^3r \int d^3r' \frac{e^2 n(\vec{r}) n(\vec{r}')}{4\pi\epsilon_0 |\vec{r} - \vec{r}'|} + \underbrace{E_{xc}[n] + T_s[n] - \text{Tr}[\frac{\hat{\rho}^2}{2m_e} \hat{\rho}]}_{U_{xc}}, \quad (1.14)$$

where  $n(r) = \langle \vec{r} | \hat{\rho} | \vec{r} \rangle$ .  $U_{xc}$  is the interaction part at full interaction strength.

**In practice we use  $E_{xc}$  instead of  $U_{xc}$**

The term  $Q_{\beta}^{DFT,\hat{W}_1}$  describes the DFT contribution when the interaction is limited to  $\hat{W}_1$ . The basic ideas behind the double-counting correction have been described in a previous paper[4].

Both, the Hartree-Fock term and the DFT double-counting term contain the same Hartree energy, which cancels exactly.

3. a **dynamic-correlation correction**

$$Q_{\beta}^{dyn,\hat{W}_2} \stackrel{\text{def}}{=} Q_{\beta}^{\hat{W}_2} - Q_{\beta}^{HF,\hat{W}_2} \quad (1.15)$$

contains an interaction term calculated on the highest level of theory such as a many-body calculation. The corresponding Hartree-Fock term is immediately subtracted as double counting correction.

The dynamic-correlation correction uses an interaction  $\hat{W}_2$ . In the spirit of the local approximation, the interaction of this term is limited to on-site terms only. The evaluation of the dynamic correlations is the only contribution that requires expensive many-particle calculations.

This term is further elaborated in section 1.3

### Stationary conditions for wave functions and occupations

Let us now explore the stationary conditions in Eq. 1.10: We use the representation of the wave functions  $|\psi_n\rangle = \sum_a |\chi_a\rangle \langle \pi_a | \psi_n \rangle$  in terms of local orbitals  $|\chi_a\rangle$ .

$$\begin{aligned} & \left[ \frac{\hat{p}^2}{2m_e} + v_{\text{ext}} + v_{\text{Hartree}} + v_{\text{xc}} + \sum_{a,b} |\pi_a\rangle \left( \frac{\delta Q_{\text{HF},\beta}^{\hat{W}_1}}{\delta \rho_{a,b}} - \frac{\delta Q_{\text{DFT},\beta}^{\hat{W}_1}}{\delta \rho_{a,b}} + \frac{\delta Q_{\text{dyn},\beta}^{\hat{W}_2}}{\delta \rho_{a,b}} \right) \langle \pi_b| \right] |\psi_n\rangle \\ &= \sum_m |\psi_m\rangle \Lambda_{m,n} \frac{1}{f_n} \end{aligned} \quad (1.16)$$

$$\begin{aligned} \langle \psi_n | & \left[ \frac{\hat{p}^2}{2m_e} + v_{\text{ext}} + v_{\text{Hartree}} + v_{\text{xc}} + \sum_{a,b} |\pi_a\rangle \left( \frac{\delta Q_{\text{HF},\beta}^{\hat{W}_1}}{\delta \rho_{a,b}} - \frac{\delta Q_{\text{DFT},\beta}^{\hat{W}_1}}{\delta \rho_{a,b}} + \frac{\delta Q_{\text{dyn},\beta}^{\hat{W}_2}}{\delta \rho_{a,b}} \right) \langle \pi_b| \right] |\psi_n\rangle \\ & - \mu + \frac{1}{\beta} \ln \left[ \frac{f_n}{1 - f_n} \right] = 0 \end{aligned} \quad (1.17)$$

These equations specify the optimum density matrix as

#### EQUILIBRIUM DENSITY MATRIX

$$\hat{\rho} = \left[ \hat{1} + e^{\beta(\hat{h}_{\text{eff}} - \mu)} \right]^{-1} \quad (1.18)$$

with the effective Hamiltonian

$$\hat{h}_{\text{eff}} = \frac{\hat{p}^2}{2m_e} + v_{\text{ext}} + v_{\text{Hartree}} + v_{\text{xc}} + \sum_{a,b} |\pi_a\rangle \left( \frac{\delta Q_{\text{HF},\beta}^{\hat{W}_1}}{\delta \rho_{a,b}} - \frac{\delta Q_{\text{DFT},\beta}^{\hat{W}_1}}{\delta \rho_{a,b}} + \frac{\delta Q_{\text{dyn},\beta}^{\hat{W}_2}}{\delta \rho_{a,b}} \right) \langle \pi_b| \quad (1.19)$$

Note, that, in equilibrium,  $\hat{h}_{\text{eff}} = \hat{h}$ . The latter defined in Eq. 1.5.

Eq. 1.18 can be verified from the stationary conditions by going into a basis of eigenstates of  $\hat{h}_{\text{eff}}$ .

The Hamiltonian  $\hat{h}_{\text{eff}}$  has, like  $\bar{h}$ , no spectral information. It simply reflects the occupations. If an orbital obtains partial fluctuations due to quantum fluctuations<sup>2</sup>,  $\hat{h}_{\text{eff}}$  places this state near the Fermi level, where the Fermi distribution give it the desired occupation.

The spectral information is contained in the Hamiltonian

$$\hat{h}_{\text{bare}} = \frac{\hat{p}^2}{2m_e} + v_{\text{ext}} + v_{\text{Hartree}} + v_{\text{xc}} \quad (1.20)$$

which reflects the Kohn-Sham energy levels. When the Green's function is calculated we will see that this Hamiltonian, respectively  $\mathbf{h} - \mathbf{\Gamma}$  acts as “non-interacting” Hamiltonian to which the self energy (including double counting) is added to give the quasi-particle spectrum.

<sup>2</sup>With quantum fluctuations we mean for example electron-hole pairs, which screen the bare interaction. Electron-hole pairs lead fractional occupations beyond the effect of thermal fluctuations.



### 1.3 Dynamic-correlation correction

The dynamical term  $Q_{\text{dyn},\beta}^{\hat{W}_2}$  is simply the difference of the complete term minus the Hartree-Fock contribution.

$$\begin{aligned}
 Q_{\text{dyn},\beta}^{\hat{W}_2}[\rho] &= \tilde{F}_{\beta}^{\hat{W}_2}[\rho] - \frac{1}{\beta} \text{Tr}[\rho \ln(\rho) + (\mathbf{1} - \rho) \ln(\mathbf{1} - \rho)] - \frac{1}{2} \sum_{a,b,c,d} U_{a,b,d,c}^{\hat{W}_2} [\rho_{d,a} \rho_{c,b} - \rho_{c,a} \rho_{d,b}] \\
 &= \text{stat}_{\mathbf{h}'_{\text{dyn}}} \text{stat}_{\mathbf{G}, \mathbf{\Sigma}_{\text{dyn}}} \left\{ \Phi_{\beta}^{LW}[\mathbf{G}, \hat{W}_2] - \Phi_{\beta}^{HF,LW}[\mathbf{G}, \hat{W}_2] \right. \\
 &\quad - \frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \ln \left[ \mathbf{1} - \left( (i\hbar\omega_{\nu} + \mu) \mathbf{1} - \bar{\mathbf{h}} \right)^{-1} \left( \mathbf{h}'_{\text{dyn}} + \mathbf{\Sigma}_{\text{dyn}}(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \right] \right. \\
 &\quad \left. \left. + (\mathbf{h}'_{\text{dyn}} + \mathbf{\Sigma}_{\text{dyn}}(i\omega_{\nu}) - \bar{\mathbf{h}}) \mathbf{G}(i\omega_{\nu}) - \left[ \mathbf{G}(i\omega_{\nu}) - \left( (i\hbar\omega_{\nu} + \mu) \mathbf{1} - \bar{\mathbf{h}} \right)^{-1} \right] \left( \mathbf{h}'_{\text{dyn}} - \bar{\mathbf{h}} \right) \right\} \right\} \\
 &\quad (1.21)
 \end{aligned}$$

We introduce a new variable

$$\mathbf{\Gamma} \stackrel{\text{def}}{=} \bar{\mathbf{h}} - \mathbf{h}'_{\text{dyn}} \quad (1.22)$$

so that

DYNAMIC CORRELATION CORRECTION

$$\begin{aligned}
 Q_{\text{dyn},\beta}^{\hat{W}_2}[\rho] &= \text{stat}_{\mathbf{\Gamma}} \text{stat}_{\mathbf{G}, \mathbf{\Sigma}_{\text{dyn}}} \left\{ \Phi_{\beta}^{LW}[\mathbf{G}, \hat{W}_2] - \Phi_{\beta}^{HF,LW}[\mathbf{G}, \hat{W}_2] \right. \\
 &\quad - \frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \ln \left[ \mathbf{1} - \left( (i\hbar\omega_{\nu} + \mu) \mathbf{1} - \bar{\mathbf{h}} \right)^{-1} \left( \mathbf{\Sigma}_{\text{dyn}}(i\omega_{\nu}) - \mathbf{\Gamma} \right) \right] \right. \\
 &\quad \left. \left. + (\mathbf{\Sigma}_{\text{dyn}}(i\omega_{\nu}) - \mathbf{\Gamma}) \mathbf{G}(i\omega_{\nu}) + \left[ \mathbf{G}(i\omega_{\nu}) - \left( (i\hbar\omega_{\nu} + \mu) \mathbf{1} - \bar{\mathbf{h}} \right)^{-1} \right] \mathbf{\Gamma} \right\} \right\} \\
 &\quad (1.23)
 \end{aligned}$$

If the Luttinger-Ward functional is expressed by Feynman diagrams, we can simply avoid the Hartree and the exchange term instead of subtracting the exchange contribution externally.

According to appendix D, the Hartree-Fock contribution can also be moved out of the variational part, where it can be calculated directly from the density matrix. In that case, however, the values of  $\mathbf{h}'_{\text{dyn}}$  and  $\mathbf{\Sigma}_{\text{dyn}}$  have different values than if they are kept inside the variational part. It is beneficial to keep it inside the variational principle, because the values of the self energy and the Lagrange multipliers will be smaller.

#### Derivative of the dynamical correlation

Let us form the derivative of  $Q_{\text{dyn},\beta}^{\hat{W}_2}[\rho]$  in all detail to avoid using inappropriate assumptions.

In order to simplify the discussion we introduce the non-interacting density matrix, which produces a given density matrix as

$$\bar{\mathbf{G}}(i\omega_\nu) = \left[ (i\hbar\omega_\nu + \mu)\mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \quad (1.24)$$

with  $\bar{\mathbf{h}}$  defined in Eq. 1.5 as a functional of the density matrix.

During the derivation we use the rule for derivatives of an inverse of a matrix<sup>3</sup>  $d(\bar{\mathbf{G}}^{-1}) = +\bar{\mathbf{G}}d\bar{\mathbf{h}}\bar{\mathbf{G}}$ .

We form the first variation with respect to the density matrix, keeping  $\Gamma$ ,  $\Sigma_{dyn}$  and  $\mathbf{G}$  fixed.

$$\begin{aligned} dQ_{\beta}^{dyn, \hat{W}_2} &= -\frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \left( -\frac{\beta \delta \Phi_{\beta}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} + \frac{\beta \delta \Phi_{\beta, HF}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} \right) d\mathbf{G} \right. \\ &\quad + \left[ \mathbf{1} - \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \right]^{-1} \left( -d\bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) - \bar{\mathbf{G}}(d\Sigma_{dyn} - d\Gamma) \right) \\ &\quad + \left( d\Sigma_{dyn} - \underbrace{d\Gamma}_{(A)} \right) \mathbf{G} + (\Sigma_{dyn} - \underbrace{\Gamma}_{(B)}) d\mathbf{G} + \left( \underbrace{d\mathbf{G}}_{(B)} - d\bar{\mathbf{G}} \right) \Gamma + \left( \underbrace{\mathbf{G}}_{(A)} - \bar{\mathbf{G}} \right) d\Gamma \left. \right\} \\ &= -\frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \left[ \Sigma_{dyn} - \frac{\beta \delta \Phi_{\beta}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} + \frac{\beta \delta \Phi_{\beta, HF}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} \right] d\mathbf{G} \right. \\ &\quad + \left[ -\left[ \mathbf{1} - \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \right]^{-1} \bar{\mathbf{G}} + \mathbf{G} \right] d\Sigma_{dyn} \\ &\quad + \left[ \left[ \mathbf{1} - \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \right]^{-1} \bar{\mathbf{G}} - \bar{\mathbf{G}} \right] d\Gamma \\ &\quad - \left[ \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \left[ \mathbf{1} - \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \right]^{-1} \bar{\mathbf{G}} \right] d\bar{\mathbf{h}} \\ &\quad \left. - d\bar{\mathbf{G}}\Gamma \right\} \\ &= -\frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \left[ \Sigma_{dyn} - \frac{\beta \delta \Phi_{\beta}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} + \frac{\beta \delta \Phi_{\beta, HF}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} \right] d\mathbf{G} \right. \\ &\quad + \left[ \mathbf{G} - \left[ \mathbf{1} - \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \right]^{-1} \bar{\mathbf{G}} \right] d\Sigma_{dyn} \\ &\quad + \left[ \left[ \mathbf{1} - \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \right]^{-1} \bar{\mathbf{G}} - \bar{\mathbf{G}} \right] (d\Gamma - d\bar{\mathbf{h}}) \\ &\quad + \underbrace{\left[ \left[ \mathbf{1} - \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \right]^{-1} \bar{\mathbf{G}} - \bar{\mathbf{G}} - \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \left[ \mathbf{1} - \bar{\mathbf{G}}(\Sigma_{dyn} - \Gamma) \right]^{-1} \bar{\mathbf{G}} \right]}_{=0} d\bar{\mathbf{h}} \\ &\quad \left. - d\bar{\mathbf{G}}\Gamma \right\} \end{aligned} \quad (1.26)$$

$$\mathbf{A}\mathbf{A}^{-1} = \mathbf{1} \quad \Rightarrow \quad (d\mathbf{A})\mathbf{A}^{-1} + \mathbf{A}(d\mathbf{A}^{-1}) = 0 \quad \Rightarrow \quad d\mathbf{A}^{-1} = -\mathbf{A}^{-1}(d\mathbf{A})\mathbf{A}^{-1} \quad (1.25)$$

Thus we obtain

$$\begin{aligned}
dQ_{\beta}^{dyn, \hat{W}_2} = & -\frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \left[ \Sigma_{dyn}(i\omega_{\nu}) - \frac{\beta \delta \Phi_{\beta}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} + \frac{\beta \delta \Phi_{\beta, HF}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} \right] d\mathbf{G}(i\omega_{\nu}) \right\} \\
& -\frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \left[ \mathbf{G}(i\omega_{\nu}) - \left[ \mathbf{1} - \bar{\mathbf{G}}(i\omega_{\nu}) \left( \Sigma_{dyn}(i\omega_{\nu}) - \mathbf{\Gamma} \right) \right]^{-1} \bar{\mathbf{G}}(i\omega_{\nu}) \right] d\Sigma_{dyn}(i\omega_{\nu}) \right\} \\
& -\frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \left[ \mathbf{1} - \bar{\mathbf{G}} \left( \Sigma_{dyn} - \mathbf{\Gamma} \right) \right]^{-1} \bar{\mathbf{G}} - \bar{\mathbf{G}} \right\} (d\mathbf{\Gamma} - d\bar{\mathbf{h}}) \\
& + \text{Tr} [\mathbf{\Gamma} d\rho]
\end{aligned} \tag{1.27}$$

Thus the stationary conditions are

- the definition of the dynamic self energy

$$\Sigma_{dyn}(i\omega_{\nu}) = \frac{\beta \delta \Phi_{\beta}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} - \frac{\beta \delta \Phi_{\beta, HF}^{LW}}{\delta \mathbf{G}(i\omega_{\nu})} \tag{1.28}$$

- Dyson's equation defining the Green's function

$$\mathbf{G}(i\omega_{\nu}) = \left[ \mathbf{1} - \bar{\mathbf{G}}(i\omega_{\nu}) \left( \Sigma_{dyn}(i\omega_{\nu}) - \mathbf{\Gamma} \right) \right]^{-1} \bar{\mathbf{G}}(i\omega_{\nu}) \tag{1.29}$$

- and the density-matrix constraint

$$\frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \left[ \mathbf{1} - \bar{\mathbf{G}} \left( \Sigma_{dyn} - \mathbf{\Gamma} \right) \right]^{-1} \bar{\mathbf{G}} - \bar{\mathbf{G}} \right\} \tag{1.30}$$

If the stationary conditions are obeyed, the derivative with respect to the density matrix is simply the matrix  $\mathbf{\Gamma}$ .

#### DENSITY-MATRIX DERIVATIVE OF THE DYNAMIC CORRELATION CORRECTION

$$\frac{\delta Q_{\beta}^{dyn, \hat{W}_2}[\rho]}{\delta \rho} = \mathbf{\Gamma} \tag{1.31}$$

The matrix  $\mathbf{\Gamma}$  replaces  $\mathbf{h}'$  as Lagrange parameter for the density matrix constraint. It vanishes in the Hartree-Fock approximation as seen from BPP-Eq.45-46[1]. Instead of the Hartree-Fock self energy  $\Sigma^{HF}$ , we could have taken any other static self energy.

One may ask, where the self energy contribution from static correlations is, because both,  $\Sigma$  and  $\Sigma^{HF}$  mentioned here refer to the second interaction using  $\hat{W}_2$ . The Terms related to  $\hat{W}_1$  are contained in the density matrix and  $\bar{\mathbf{h}}$ .

## 1.4 Local orbitals

### 1.4.1 Projection onto local orbitals

In order to integrate DMFT into the DFT code, we define first a local basis set of orbitals  $|\chi_a\rangle$ . These orbitals are not orthogonal. The orbitals are spin orbitals, that is, each is a two-component wave function with a spin-up and a spin-down component. If the spin orbitals are eigenstates of  $\hat{S}_z$ , one or the other of the components vanishes.

The decomposition of the Kohn-Sham wave functions, which in the the context of rDMFT are the natural orbitals, is obtained via the projector functions  $\langle\pi_a|$  as

$$|\psi_n\rangle = \sum_a |\chi_a\rangle \langle\pi_a|\psi_n\rangle + |\delta\psi_n\rangle \quad (1.32)$$

where the projector functions obey the bi-orthogonality condition

$$\langle\pi_a|\chi_b\rangle = \delta_{a,b}, \quad (1.33)$$

and where  $|\delta\psi_n\rangle$  is a remainder which is left over if the local orbitals do not form a complete basis set. This remainder has the property

$$\langle\pi_a|\delta\psi_n\rangle = 0 \quad (1.34)$$

However, there is no orthogonality between the orbitals and the remainder, that is  $\langle\chi_\alpha|\delta\psi_n\rangle \neq 0$ .

In order to estimate the importance of the remainder, consider the following thought experiment. Perform calculation with the self-consistent potential and introduce the constraints  $\langle\pi_\alpha|\bar{\psi}_n\rangle = 0$  for all  $\alpha$ . The constraints will push the wave functions up in energy. These wave functions serve as basis functions for the remainder  $|\delta\psi_n\rangle = \sum_m |\bar{\psi}_m\rangle c_{m,n}$ . If there are many constraints and there are no larger gaps between the atoms, the constrained wave functions  $|\bar{\psi}_m\rangle$  will have high energies and therefore we can assume that they do not contribute much to the wave functions. As conclusion from our thought experiment comes the notion that we can assume the remainder to be unimportant, if there are many projectors and if they do not leave any large gaps in space.

### 1.4.2 Energy correction

Here we consider changes of the correlation corrections when we move from a Bra-Ket notation to a matrix notation.

$$Q_\beta^{\hat{W}}[\hat{\rho}] = Q_\beta^{\hat{W}} \left[ \sum_n |\psi_n\rangle f_n \langle\psi_n| \right] = Q_\beta^{\hat{W}} \left[ \sum_{a,b} |\chi_a\rangle \underbrace{\left( \sum_n \langle\pi_a|\psi_n\rangle f_n \langle\psi_n|\pi_b\rangle \right)}_{\rho_{a,b}} \langle\chi_b| \right] \quad (1.35)$$

Below we consider also the energy contribution as functional of the density matrix in matrix form

$$Q_\beta^{\hat{W}}[\boldsymbol{\rho}] \stackrel{\text{def}}{=} Q_\beta^{\hat{W}} \left[ \sum_{a,b} |\chi_a\rangle \rho_{a,b} \langle\chi_b| \right] \quad (1.36)$$

The derivatives are

$$\begin{aligned}\frac{\delta Q_{\beta}^{\hat{W}}}{\delta \langle \psi_n |} \frac{1}{f_n} &= \sum_{a,b} |\pi_b\rangle \langle \chi_b| \frac{\delta Q_{\beta}^{\hat{W}}}{\delta \hat{\rho}} |\chi_a\rangle \langle \pi_a | \psi_n \rangle \\ \frac{\delta Q_{\beta}^{\hat{W}}}{\delta f_n} &= \sum_{a,b} \langle \psi_n | \pi_b \rangle \langle \chi_b| \frac{\delta Q_{\beta}^{\hat{W}}}{\delta \hat{\rho}} |\chi_a\rangle \langle \pi_a | \psi_n \rangle\end{aligned}\quad (1.37)$$

In the following we introduce a new symbol, namely

$$\frac{\delta Q_{\beta}^{\hat{W}}}{\delta \rho_{b,a}} \stackrel{\text{def}}{=} \langle \chi_b | \frac{\delta Q_{\beta}^{\hat{W}}}{\delta \hat{\rho}} | \chi_a \rangle \quad (1.38)$$

### 1.4.3 U-tensor

The interaction  $\hat{W}$

$$\hat{W} = \frac{1}{2} \sum_{a,b,c,d} U_{a,b,d,c} \hat{c}_a^{\dagger} \hat{c}_b^{\dagger} \hat{c}_c \hat{c}_d \quad (1.39)$$

is expressed by the **U-tensor**

$$U_{a,b,c,d} = \sum_{\sigma, \sigma'} \int d^3 r \int d^3 r' \frac{e^2 \chi_a^*(\vec{r}, \sigma) \chi_b^*(\vec{r}', \sigma') \chi_c(\vec{r}, \sigma) \chi_d(\vec{r}', \sigma')}{4\pi\epsilon_0 |\vec{r} - \vec{r}'|} \quad (1.40)$$

The U-tensor can then be approximated to yield  $\hat{W}_1$  and  $\hat{W}_2$ . Typically, these approximations amount to multiplying matrix elements with scale factors and to leaving certain elements out completely.

### 1.4.4 Non-orthonormal orbitals

The main difference[4] between orthonormal and non-orthonormal basisets is that the commutator relation

$$\left[ \hat{c}_a^{\dagger}, \hat{c}_b \right]_{+} = \langle \pi_b | \pi_a \rangle \quad (1.41)$$

### Feynman diagrams

A Feynmandiagram in terms of orthonormal orbitals can be transformed easily into any other one-particle basis, by introducing a unit matrix in the form  $\mathbf{AA}^{-1}$  in between every Green's function and every interaction.

Let the original basis be  $\{|\psi_n\rangle\}$  and a new basis  $\{|\chi_a\rangle\}$  so that

$$|\psi_n\rangle = \sum_a |\chi_a\rangle \langle \pi_a | \psi_n \rangle \quad (1.42)$$

Then the Green's function has the form

$$\hat{G} = \sum_{m,n} |\psi_m\rangle \langle \psi_m | \hat{G} | \psi_n \rangle \langle \psi_n | = \sum_{a,b} |\chi_a\rangle \underbrace{\langle \pi_a | \hat{G} | \pi_b \rangle}_{G_{a,b}} \langle \chi_b | \quad (1.43)$$

Thus if the U-tensor is evaluated using the orbitals  $|\chi\rangle$ , it can directly be used with the Green's function  $\langle\pi_a|\hat{G}|\pi_b\rangle$ . The final result will be invariant if all sums are traced out.

#### 1.4.5 Lattice Green's functions in a finite, non-orthonormal basis

The Green's function results as inverse of a hamiltonian. The Green's function we need should only act in a sub-Hilbert space of correlated orbitals. Therefore we introduce a projection operator

$$\hat{P}_\chi = \sum_{a,b} |\chi_a\rangle S_{a,b}^{-1} \langle\chi_b| \quad \text{with} \quad \sum_c S_{a,c}^{-1} \langle\chi_c|\chi_b\rangle = \delta_{a,b} \quad (1.44)$$

onto the correlated orbitals. The projection operator is idempotent, i.e.  $\hat{P}_\chi^2 = \hat{P}_\chi$

The Green's function is that operator that acts in the Hilbert space of the correlated orbitals, and that inverts the hamiltonian projected onto that hilbert space. Thus the following equation defines the greens function.

$$\hat{G}\hat{P}_\chi \left[ (i\hbar\omega_\nu + \mu)\hat{1} - \hat{h} - \sum_{c,d} |\pi_c\rangle \Sigma_{c,d}(i\omega_\nu) \langle\pi_d| \right] \hat{P}_\chi = \hat{P}_\chi \quad (1.45)$$

$$\begin{aligned} \underbrace{|\chi\rangle \mathbf{G} \langle\chi|}_{\hat{G}} \underbrace{|\chi\rangle \mathbf{S}^{-1} \langle\chi|}_{\hat{P}_\chi} \left[ (i\hbar\omega_\nu + \mu)\hat{1} - \hat{h} - \sum_{c,d} |\pi_c\rangle \Sigma_{c,d}(i\omega_\nu) \langle\pi_d| \right] \underbrace{|\chi\rangle \mathbf{S}^{-1} \langle\chi|}_{\hat{P}_\chi} &= \underbrace{|\chi\rangle \mathbf{S}^{-1} \langle\chi|}_{\hat{P}_\chi} \\ \mathbf{G} \langle\chi| \left[ (i\hbar\omega_\nu + \mu)\hat{1} - \hat{h} - \sum_{c,d} |\pi_c\rangle \Sigma_{c,d}(i\omega_\nu) \langle\pi_d| \right] |\chi\rangle &= \mathbf{1} \\ \mathbf{G} \left[ (i\hbar\omega_\nu + \mu)\mathbf{S} - \langle\chi|\hat{h}|\chi\rangle - \Sigma(i\omega_\nu) \right] &= \mathbf{1} \\ \mathbf{G} \left[ (i\hbar\omega_\nu + \mu)\mathbf{S} - \sum_{n=1}^{\infty} \langle\chi|\psi_n\rangle \epsilon_n \langle\psi_n|\chi\rangle - \Sigma(i\omega_\nu) \right] &= \mathbf{1} \\ \mathbf{G} \left[ (i\hbar\omega_\nu + \mu)\mathbf{S} - \epsilon_x \mathbf{S} - \sum_{n=1}^{\infty} \langle\chi|\psi_n\rangle (\epsilon_n - \epsilon_x) \langle\psi_n|\chi\rangle - \Sigma(i\omega_\nu) \right] &= \mathbf{1} \end{aligned} \quad (1.46)$$

Thus we obtain

$$\mathbf{G} = \left[ (i\hbar\omega_\nu + \mu)\mathbf{S} - \epsilon_x \mathbf{S} - \sum_{n=1}^{\infty} \langle\chi|\psi_n\rangle (\epsilon_n - \epsilon_x) \langle\psi_n|\chi\rangle - \Sigma(i\omega_\nu) \right]^{-1} \quad (1.47)$$

Now we modify the Hamiltonian by changing the spectrum for the higher bands

$$\hat{h} \rightarrow \sum_{n=1}^{N_x} |\psi_n\rangle \epsilon_n \langle\psi_n| + \sum_{n=N_x+1}^{\infty} |\psi_n\rangle \epsilon_x \langle\psi_n| \quad (1.48)$$

This approximation allows one to limit the sum over states to a finite sum.

$$\mathbf{G} = \left[ (i\hbar\omega_\nu + \mu - \epsilon_x)\mathbf{S} - \sum_{n=1}^{N_x} \langle\chi|\psi_n\rangle (\epsilon_n - \epsilon_x) \langle\psi_n|\chi\rangle - \Sigma(i\omega_\nu) \right]^{-1} \quad (1.49)$$

After limiting the sum over the states that can be well represented by the correlated orbitals, we use the approximation  $|\psi_n\rangle = \sum_a |\chi_a\rangle \langle \pi_a | \psi_n \rangle$ , which yields

$$\mathbf{G} = \left[ (i\hbar\omega_\nu + \mu - \epsilon_x) \mathbf{S} - \sum_{n=1}^{N_x} \mathbf{S} \langle \pi | \psi_n \rangle (\epsilon_n - \epsilon_x) \langle \psi_n | \pi \rangle \mathbf{S} - \boldsymbol{\Sigma}(i\omega_\nu) \right]^{-1} \quad (1.50)$$

This matrix corresponds to

$$\mathbf{G} = \left[ (i\hbar\omega_\nu + \mu) \mathbf{S} - \bar{\mathbf{h}} + \boldsymbol{\Gamma} - \boldsymbol{\Sigma}(i\omega_\nu) \right]^{-1} \quad (1.51)$$

which allows one to make the identification (for full selfconsistency)

$$\begin{aligned} \epsilon_x \left[ \mathbf{S} - \sum_{n=1}^{N_x} \mathbf{S} \langle \pi | \psi_n \rangle \langle \psi_n | \pi \rangle \mathbf{S} \right] - \underbrace{\sum_{n=1}^{N_x} \mathbf{S} \langle \pi | \psi_n \rangle \epsilon_n \langle \psi_n | \pi \rangle \mathbf{S}}_{\rightarrow \frac{\hat{p}^2}{2m} + V_{\text{ext}} + V_H + V_{xc}} &= -\bar{\mathbf{h}} + \boldsymbol{\Gamma} \\ \Rightarrow \quad \bar{\mathbf{h}} = \underbrace{\sum_{n=1}^{N_x} \mathbf{S} \langle \pi | \psi_n \rangle \epsilon_n \langle \psi_n | \pi \rangle \mathbf{S}}_{\rightarrow \frac{\hat{p}^2}{2m} + V_{\text{ext}} + V_H + V_{xc}} - \epsilon_x \left[ \mathbf{S} - \sum_{n=1}^{N_x} \mathbf{S} \langle \pi | \psi_n \rangle \langle \psi_n | \pi \rangle \mathbf{S} \right] + \boldsymbol{\Gamma} \end{aligned} \quad (1.52)$$

## Projections

In the last step of the previous section we did the step

$$\sum_{n=1}^{N_x} \langle \chi | \psi_n \rangle \langle \psi_n | \chi \rangle \approx \sum_{n=1}^{N_x} \mathbf{S} \langle \pi | \psi_n \rangle \langle \psi_n | \pi \rangle \mathbf{S} \quad (1.53)$$

which is only approximative.

An allowed step is

$$\begin{aligned} \langle \chi | \psi_n \rangle \langle \psi_n | \chi \rangle &= \langle \chi | \chi \rangle \langle \pi | \hat{\mathcal{P}}_\chi | \psi_n \rangle \langle \psi_n | \hat{\mathcal{P}}_\chi | \pi \rangle \langle \chi | \chi \rangle = \mathbf{S} \langle \pi | \hat{\mathcal{P}}_\chi | \psi_n \rangle \langle \psi_n | \hat{\mathcal{P}}_\chi | \pi \rangle \mathbf{S} \\ &\neq \mathbf{S} \langle \pi | \psi_n \rangle \langle \psi_n | \pi \rangle \mathbf{S} \end{aligned} \quad (1.54)$$

**This is only a sketch...**

Consider the set of wave functions obtained by projection onto the correlated orbitals.

$$\hat{\mathcal{P}}_\chi | \psi_n \rangle \quad (1.55)$$

We obtain a strict identity

$$\hat{\mathcal{P}}_\chi | \psi_n \rangle = \sum_{\alpha} |\chi_\alpha\rangle \langle \pi_\alpha | \hat{\mathcal{P}}_\chi | \psi_n \rangle \quad (1.56)$$

However, the wave functions  $\hat{\mathcal{P}}_\chi | \psi_n \rangle$  are no more orthonormal, and  $\langle \pi_\alpha | \hat{\mathcal{P}}_\chi | \psi_n \rangle \neq \langle \pi_\alpha | \hat{\mathcal{P}}_\chi | \psi_n \rangle$ .

### Calculate $\bar{h}$ from the density matrix

The constraint is

$$\sum_{\nu} e^{i\hbar\omega_{\nu}\beta 0^+} G_{a,b}(i\omega_{\nu}) = \rho_{a,b} (= \sum_n \langle \pi_a | \psi_n \rangle f_n \langle \psi_n | \pi_b \rangle) \quad (1.57)$$

Thus we look for a hamiltonian  $\bar{h}$  whose non-interacting Greens function satisfies the constraint

$$\rho = \sum_{\nu} e^{i\hbar\omega_{\nu}\beta 0^+} \left[ (i\hbar\omega_{\nu} + \mu) \mathbf{S} - \bar{h} \right]^{-1} \quad (1.58)$$

$$\Rightarrow \quad \sqrt{\mathbf{S}} \rho \sqrt{\mathbf{S}} = \sum_{\nu} e^{i\hbar\omega_{\nu}\beta 0^+} \left[ (i\hbar\omega_{\nu} + \mu) \mathbf{1} - (\mathbf{S})^{-\frac{1}{2}} \bar{h} (\mathbf{S})^{-\frac{1}{2}} \right]^{-1} \quad (1.59)$$

In the next step we diagonalize  $\sqrt{\mathbf{S}} \rho \sqrt{\mathbf{S}}$

$$\sqrt{\mathbf{S}} \rho \sqrt{\mathbf{S}} \mathbf{U} = \mathbf{U} \bar{f} \quad (1.60)$$

which yields effective occupations  $\bar{f}$ . The eigenvalue equation is equivalent with the generalized eigenvalue problem

$$\rho(\sqrt{\mathbf{S}} \mathbf{U}) = \mathbf{S}^{-1} (\sqrt{\mathbf{S}} \mathbf{U}) \bar{f} \quad (1.61)$$

Now we calculate  $\bar{\epsilon}$  so that

$$\bar{f} = \left[ 1 + e^{\beta(\bar{\epsilon} - \mu \mathbf{1})} \right]^{-1} \quad \Leftrightarrow \quad \bar{\epsilon} = \mu \mathbf{1} + \frac{1}{\beta} \ln \left( \frac{\mathbf{1} - \bar{f}}{\bar{f}} \right) \quad (1.62)$$

Then we can represent  $\bar{f}$  by the Matsubara sum

$$\bar{f} = \sum_{\nu} e^{i\hbar\omega_{\nu}\beta 0^+} \left[ (i\hbar\omega_{\nu} + \mu) \mathbf{1} - \bar{\epsilon} \right]^{-1} \quad (1.63)$$

Thus we obtain

$$\rho = \mathbf{S}^{-\frac{1}{2}} \mathbf{U} \bar{f} \mathbf{U}^{\dagger} \mathbf{S}^{-\frac{1}{2}} = \sum_{\nu} e^{i\hbar\omega_{\nu}\beta 0^+} \mathbf{S}^{-\frac{1}{2}} \mathbf{U} \left[ (i\hbar\omega_{\nu} + \mu) \mathbf{1} - \bar{\epsilon} \right]^{-1} \mathbf{U}^{\dagger} \mathbf{S}^{-\frac{1}{2}} \quad (1.64)$$

This yields

$$\begin{aligned} \bar{\mathbf{G}} &= \mathbf{S}^{-\frac{1}{2}} \mathbf{U} \left[ (i\hbar\omega_{\nu} + \mu) \mathbf{1} - \bar{\epsilon} \right]^{-1} \mathbf{U}^{\dagger} \mathbf{S}^{-\frac{1}{2}} \\ &= \left[ (i\hbar\omega_{\nu} + \mu) \mathbf{S} - \underbrace{\sqrt{\mathbf{S}} \mathbf{U} \bar{\epsilon} \mathbf{U}^{\dagger} \sqrt{\mathbf{S}}}_{\bar{h}} \right]^{-1} = \left[ (i\hbar\omega_{\nu} + \mu) \mathbf{S} - \underbrace{\mathbf{V} \bar{\epsilon} \mathbf{V}^{\dagger}}_{\bar{h}} \right]^{-1} \end{aligned} \quad (1.65)$$

Comparison yields

$$\bar{h} = \mathbf{V} \left[ \mu \mathbf{1} + \frac{1}{\beta} \ln \left( \frac{\mathbf{1} - \bar{f}}{\bar{f}} \right) \right] \mathbf{V}^{\dagger} \quad (1.66)$$

where

$$\rho \mathbf{V} = \mathbf{S}^{-1} \mathbf{V} \bar{f} \quad \text{with} \quad \mathbf{V}^{\dagger} \mathbf{S}^{-1} \mathbf{V} = \mathbf{1} \quad (1.67)$$



### 1.4.6 Incompleteness

One difficulty stems from the fact that the natural orbitals, the local orbitals and the projector functions span different spaces.

We choose a solution by introducing a new type of algebra, that is designed to give certain sum rules exactly. In this algebra we consider the states

$$|\kappa_n\rangle \stackrel{\text{def}}{=} \sum_a |\chi_a\rangle \langle \pi_a | \psi_n \rangle \quad (1.68)$$

as complete orthonormal basisset.

Thus,

$$\delta_{m,n} = \langle \kappa_m | \kappa_n \rangle = \sum_{a,b} \langle \psi_m | \pi_a \rangle \langle \chi_a | \chi_b \rangle \langle \pi_b | \psi_n \rangle \quad (1.69)$$

and

$$\hat{1} = \sum_{a,b,n} |\chi_a\rangle \langle \pi_a | \psi_n \rangle \langle \psi_n | \pi_b \rangle \langle \chi_b | \quad (1.70)$$

I introduce

$$S_{a,b}^{-1} \stackrel{\text{def}}{=} \sum_n \langle \pi_a | \psi_n \rangle \langle \psi_n | \pi_b \rangle \quad (1.71)$$

and redefine the overlap  $\langle \chi_a | \chi_b \rangle$  as the inverse of  $\mathbf{S}$ . We could consider this as a redefinition of the scalar product between the states  $|\kappa_n\rangle$ .

### Derivatives

If the incompleteness is taken into account in the way suggested above, the density matrix functional not only depends on the density matrix, but also on the inverse overlap matrix.

$$F[\rho, \mathbf{S}^{-1}] \quad (1.72)$$

The derivatives with respect to the natural orbitals and the occupations are

$$\begin{aligned} dF &= \frac{\partial F}{\partial \rho} d\rho + \frac{\partial F}{\partial \mathbf{S}^{-1}} d\mathbf{S}^{-1} \\ &= \sum_{\alpha,\beta} \frac{\partial F}{\partial \rho_{\beta,\alpha}} \left( \sum_n \langle \pi_\beta | \delta \psi_n \rangle f_n \langle \psi_n | \pi_\alpha \rangle + \sum_n \langle \pi_\beta | \psi_n \rangle df_n \langle \psi_n | \pi_\alpha \rangle + \sum_n \langle \pi_\beta | \psi_n \rangle f_n \langle \delta \psi_n | \pi_\alpha \rangle \right) \\ &\quad + \frac{\partial F}{\partial \mathbf{S}_{\beta,\alpha}^{-1}} \left( \sum_n \langle \pi_\beta | \delta \psi_n \rangle \langle \psi_n | \pi_\alpha \rangle + \sum_n \langle \pi_\beta | \psi_n \rangle \langle \delta \psi_n | \pi_\alpha \rangle \right) \\ &= \sum_n df_n \left( \sum_{\alpha,\beta} \langle \psi_n | \pi_\alpha \rangle \frac{\partial F}{\partial \rho_{\beta,\alpha}} \langle \pi_\beta | \psi_n \rangle \right) \\ &\quad + \sum_n f_n \sum_{\alpha,\beta} \langle \delta \psi_n | \pi_\alpha \rangle \left[ \frac{\partial F}{\partial \rho_{\beta,\alpha}} \langle \pi_\beta | \psi_n \rangle + \frac{\partial F}{\partial \mathbf{S}_{\beta,\alpha}^{-1}} \langle \pi_\beta | \psi_n \rangle \frac{1}{f_n} \right] + c.c. \end{aligned} \quad (1.73)$$

Problematic is the term  $\frac{1}{f_n}$ , which diverges for empty states.

For fully occupied states the two terms, those containing the derivative with respect to the density matrix and those with respect to the inverse overlap, act simultaneously. I believe that the two terms together suppress changes in the total energy that are due that the natural orbitals overlap better with the correlated orbitals.

Let us single out the term that acts on empty or partially occupied states.

$$\begin{aligned}
 dF = & \sum_n df_n \sum_{\alpha,\beta} \left[ \langle \psi_n | \pi_\alpha \rangle \left( \frac{\partial F}{\partial \rho_{\beta,\alpha}} + \frac{\partial F}{\partial \mathbf{S}_{\beta,\alpha}^{-1}} \right) \langle \pi_\beta | \psi_n \rangle + \langle \psi_n | \pi_\alpha \rangle \frac{\partial F}{\partial \mathbf{S}_{\beta,\alpha}^{-1}} \langle \pi_\beta | \psi_n \rangle \frac{1-f_n}{f_n} \right] \\
 & + \underbrace{\sum_n df_n \sum_{\alpha,\beta} \left[ -\langle \psi_n | \pi_\alpha \rangle \frac{\partial F}{\partial \mathbf{S}_{\beta,\alpha}^{-1}} \langle \pi_\beta | \psi_n \rangle - \langle \psi_n | \pi_\alpha \rangle \frac{\partial F}{\partial \mathbf{S}_{\beta,\alpha}^{-1}} \langle \pi_\beta | \psi_n \rangle \frac{1-f_n}{f_n} \right]}_{\text{correction}} \\
 & + \sum_n f_n \sum_{\alpha,\beta} \langle \delta \psi_n | \pi_\alpha \rangle \left[ \left( \frac{\partial F}{\partial \rho_{\beta,\alpha}} + \frac{\partial F}{\partial \mathbf{S}_{\beta,\alpha}^{-1}} \right) \langle \pi_\beta | \psi_n \rangle + \frac{\partial F}{\partial \mathbf{S}_{\beta,\alpha}^{-1}} \langle \pi_\beta | \psi_n \rangle \frac{1-f_n}{f_n} \right] + c.c.
 \end{aligned} \tag{1.74}$$

The term in the last line is a contribution to the force acting on the wave functions. The kohn-sham energies  $\epsilon_n$  are usually determined from this force acting on the wave functions. In the present case, however, we need to add the term marked “correction” in the second line to the Kohn-Sham energies.

One potential remedy of the singularity for empty states is to replace  $\frac{1-f}{f}$  by  $1-f$ . For states that are strongly occupied, that is for  $f \rightarrow 1$  the two functions are similar. For weakly occupied states, that is for  $f \rightarrow 0$  the term is there, but not divergent.

## 1.5 Variational parameters for the Green's function

### Representation of the Green's function

The Green's function is expressed in a Bloch representation.

$$\mathbf{G}_{\vec{t},\vec{t}'}(i\omega_\nu) = \frac{1}{N_k} \sum_{\vec{k}} \mathbf{G}_{\vec{k}}(i\omega_\nu) e^{i\vec{k}(\vec{t}-\vec{t}')} \tag{1.75}$$

where  $N_k$  is the number of  $\vec{k}$  points on the chosen numerical grid in the first unit cell.  $\mathbf{G}_{\vec{t},\vec{t}'}$  is a matrix with dimension equal to the number of local orbitals in a given unit cell, and it connects these orbitals in two unit cells specified by the translation vectors  $\vec{t}$  and  $\vec{t}'$ . **Check the sign in the exponential!**

The Green's function can be expressed by the values at the positive Matsubara frequencies up to a maximum value  $\omega_x$  and the first three Laurent expansion terms.

$$\mathbf{G}_{\vec{k}}(i\omega_\nu) = \begin{cases} \mathbf{G}_{\vec{k}}(i\omega_\nu) & \text{for } 0 < \omega_\nu \leq \omega_x \\ \mathbf{G}_{\vec{k}}^\dagger(-i\omega_\nu) & \text{for } -\omega_x \leq \omega_\nu < 0 \\ \sum_{j=1,2} \frac{1}{(i\hbar\omega_\nu)^j} \mathcal{G}_{\vec{k}}^{(j)} & \text{for } |\omega_\nu| > \omega_x \end{cases} \tag{1.76}$$

where  $\mathcal{G}_{\vec{k}}^{(0)} = \mathbf{S}_{\vec{k}}^{-1}$ .

### Variational parameters

However, we treat this representation as the derived form. By employing the stationary condition for the self energy, we obtain Dyson's equation, which expresses the Green's function by the self energy. Using the same stationary condition with respect to the Laurent expansion terms of the self energy, one obtains a consistent definition for the Laurent expansion terms of the Green's function.

The parameters entering in the representation Eq. 1.76 of the Greens function given above are obtained from the fundamental variables as.

$$\begin{aligned}
 G_{\vec{k}}(i\omega_\nu) &= \left[ (i\hbar\omega_\nu + \mu)S_{\vec{k}} - \bar{h}_{\vec{k}} - [\Sigma_{dyn}(i\omega_\nu) - \Gamma_{\vec{k}}] \right]^{-1} \quad \text{for } 0 < \omega_\nu \leq \omega_x \\
 \mathcal{G}_{\vec{k}}^{(1)} &\stackrel{\text{Eq. B.27}}{=} S_{\vec{k}}^{-1} \\
 \mathcal{G}_{\vec{k}}^{(2)} &\stackrel{\text{Eq. B.27}}{=} -S_{\vec{k}}^{-1} \left[ \mu S_{\vec{k}} - \bar{h}_{\vec{k}} + \Gamma_{\vec{k}} - S_{dyn}^{(0)} \right] S_{\vec{k}}^{-1}
 \end{aligned} \tag{1.77}$$

The leading order in the Laurent expansion of the Green's function is not a dynamical variable, because it does not depend on any of the variational parameters in the definition of the dynamical correlation.

Thus the self energy is treated as variable that is iteratively optimized. This optimization is guided by the derivative with respect to the Green's function.

$$\Sigma_{dyn}(+) = \Sigma_{dyn}(0) + \alpha \left[ \frac{\beta \delta \Phi^{LW,dyn}}{\delta G} \Big|_{G[\Sigma(0)]} - \Sigma_{dyn}(0) \right] \tag{1.78}$$

The fundamental variables are

- the site-diagonal and therefore k-independent dynamic self energy  $\Sigma_{dyn,R}(i\omega_\nu) = \Sigma_R(i\omega_\nu) - \Sigma_{HF,R}(i\omega_\nu)$ . The self energy for positive and negative Matsubara frequencies are related by a hermitean conjugation  $\Sigma_{dyn,R}(-i\omega_\nu) = \Sigma_{dyn,R}^\dagger(i\omega_\nu)$
- the Hamiltonian  $\bar{h}_{\vec{k}}$ , which produces the one-particle reduced density matrix as  $\rho_{\vec{k}} = \left[ \mathbf{1} + e^{\beta(\bar{h}_{\vec{k}} - \mu \mathbf{1})} \right]^{-1}$ .
- the static correction  $\Gamma_{\vec{k}}$ , that ensures that the original density matrix is obtained from the Green's function, if the dynamic self energy is taken into account.

From the site-diagonal self energies  $\Sigma_{dyn,R}$ , we obtain the global self energy  $\Sigma_{dyn}$  as direct sum. The global self energy is formally a matrix that connects each orbital in a unit cell with each other one.

$$\Sigma_{dyn}(i\omega_\nu) = \bigoplus_{R \in \text{unit cell}} \Sigma_{dyn,R}(i\omega_\nu) \tag{1.79}$$

This self energy is independent of  $\vec{k}$ .

### Local Green's function

Then local Green's function is obtained as Brillouin-zone integral

$$\mathbf{G}_R(i\omega_\nu) = \frac{1}{N_k} \sum_{\vec{k}} \mathbf{P}_R \mathbf{G}_{\vec{k}}(i\omega_\nu) \mathbf{P}_R^\dagger \quad (1.80)$$

where  $\mathbf{P}_R$  is a non-square diagonal matrix, which projects the matrix onto the subspace of orbitals on a given site.

#### 1.5.1 High-Frequency expansion of Matsubara sum

A consistent choice for the Laurent expansion for the Green's function and self energy has turned out to be non-trivial. The most consistent approach appears to write down an approximate form for the high-frequency tail of the Matsubara sum in the dynamic correlation correction, and to extract the Laurent expansion coefficients from the stationary conditions.

With the high-frequency tail, we obtain

$$\begin{aligned} & \ln[1 - \bar{\mathbf{G}}(\Sigma - \Gamma)] + (\Sigma - \Gamma)\mathbf{G} + (\mathbf{G} - \bar{\mathbf{G}})\Gamma \\ &= -\bar{\mathbf{G}}(\Sigma - \Gamma) - \sum_{n=2}^{\infty} \frac{1}{n} [\bar{\mathbf{G}}(\Sigma - \Gamma)]^n + (\Sigma - \Gamma)\mathbf{G} + (\mathbf{G} - \bar{\mathbf{G}})\Gamma \\ &= (\mathbf{G} - \bar{\mathbf{G}})(\Sigma - \Gamma) - \sum_{n=2}^{\infty} \frac{1}{n} [\bar{\mathbf{G}}(\Sigma - \Gamma)]^n + \underbrace{(\mathbf{G} - \bar{\mathbf{G}})\Gamma}_{\rightarrow \text{number constraint}} \\ &= (\mathbf{G} - \bar{\mathbf{G}})\Sigma - \sum_{n=2}^{\infty} \frac{1}{n} [\bar{\mathbf{G}}(\Sigma - \Gamma)]^n \end{aligned} \quad (1.81)$$

Side Remark:<sup>4</sup>

<sup>4</sup>In order to rationalize the form of the Kadanoff-Baym functional, we look for a functional  $X$ , that provides Dyson's equation as functional derivative of the self energy. We do not make any reference regarding the representation of the Green's function. We only require  $d\text{Tr}[A^n] = nA^{n-1}dA$  and  $\ln[1 - A] = -\sum_{n=1}^{\infty} \frac{1}{n} A^n$ .

We start with Dyson's equation

$$\mathbf{G} = \mathbf{G}_0 + \mathbf{G}_0 \Sigma \mathbf{G} = \mathbf{G}_0 + \sum_{n=1}^{\infty} \mathbf{G}_0 (\Sigma \mathbf{G}_0)^n = \sum_{n=0}^{\infty} \mathbf{G}_0 (\Sigma \mathbf{G}_0)^n \quad (1.82)$$

We define a functional  $X[G, \Sigma]$  so that

$$\left. \frac{\delta X[G, \Sigma]}{\delta \Sigma} \right|_{G_0} = \mathbf{G} - \sum_{n=0}^{\infty} \mathbf{G}_0 (\Sigma \mathbf{G}_0)^n \quad (1.83)$$

The functional has the form

$$\begin{aligned} X[G, \Sigma] &= \text{Tr} \left[ G \Sigma - \sum_{n=0}^{\infty} \frac{1}{n+1} (\Sigma \mathbf{G}_0)^{n+1} \right] = \text{Tr} \left[ G \Sigma - \sum_{n=1}^{\infty} \frac{1}{n} (\Sigma \mathbf{G}_0)^n \right] \\ &\stackrel{\text{Eq. B.17}}{=} \text{Tr} \left[ G \Sigma + \ln(1 - \Sigma \mathbf{G}_0) \right] \end{aligned} \quad (1.84)$$

The functional derivative of  $X$  with respect to the Green's function is the self energy

$$\left. \frac{\delta X[G, \Sigma]}{\delta G} \right|_{G_0} = \Sigma \quad (1.85)$$

Now we insert the Laurent expansions into Eq. 1.81

$$\begin{aligned}
& (\mathcal{G} - \bar{\mathcal{G}}) \Sigma - \sum_{n=2}^{\infty} \frac{1}{n} [\bar{\mathcal{G}} (\Sigma - \Gamma)]^n \\
&= \left[ \sum_{j=2}^{\infty} \frac{1}{(i\hbar\omega_{\nu})^j} (\mathcal{G}^{(j)} - \bar{\mathcal{G}}^{(j)}) \right] \left[ \sum_{j'=0}^{\infty} \frac{1}{(i\hbar\omega_{\nu})^{j'}} \mathcal{S}^{(j')} \right] \\
&\quad - \sum_{n=2}^{\infty} \frac{1}{n} \left[ \left( \sum_{j=1}^{\infty} \frac{1}{(i\hbar\omega_{\nu})^j} \bar{\mathcal{G}}^{(j)} \right) \left( \mathcal{S}^{(0)} - \Gamma + \sum_{j'=1}^{\infty} \frac{1}{(i\hbar\omega_{\nu})^{j'}} \mathcal{S}^{(j')} \right) \right]^n \quad (1.88)
\end{aligned}$$

So-far, no approximations have been made: Now we introduce an approximation by truncating the Laurent expansions.

- We choose to include terms in the Matsubara sum up to the second order of the inverse Matsubara frequency.
- we include only the leading order of the Laurent expansion, namely  $\mathcal{G}^{(2)}$  and  $\mathcal{S}^{(0)}$  for the self energy.

Alternatively one could include terms in the Matsubara sum up to fourth order in the inverse temperature and to include the also  $\mathcal{G}^{(3)}$  and  $\mathcal{S}^{(1)}$ . Note that the third order term in inverse temperature of the Matsubara sum does not contribute. However we will postpone this for the time being.

Another alternative would be to truncate the power-series expansion of the logarithm beyond a certain value and to include all Laurent expansion terms of the Matsubara sum which are produced by the Green's-function and self-energy expansion up to a certain order.

Then, the functional  $X$

$$\begin{aligned}
X &= -\frac{1}{\beta} \sum_{\nu, |\omega_{\nu}| > \omega_x} \text{Tr} \left\{ \frac{1}{(i\hbar\omega_{\nu})^2} \left\{ \left( \mathcal{G}^{(2)} - \bar{\mathcal{G}}^{(2)} \right) \mathcal{S}^{(0)} - \frac{1}{2} \left[ \bar{\mathcal{G}}^{(1)} \left( \mathcal{S}^{(0)} - \Gamma \right) \right]^2 \right\} \right\} \\
&\stackrel{\text{Eq. B.17}}{=} -\frac{1}{\beta} \sum_{\nu, |\omega_{\nu}| > \omega_x} \text{Tr} \left\{ \frac{1}{(i\hbar\omega_{\nu})^2} \left\{ \left( \mathcal{G}^{(2)} - \bar{\mathcal{G}}^{(2)} \right) \mathcal{S}^{(0)} - \frac{1}{2} \left[ \mathbf{s}^{-1} \left( \mathcal{S}^{(0)} - \Gamma \right) \right]^2 \right\} \right\} \quad (1.89)
\end{aligned}$$

The functional derivatives of Eq. 1.89 with respect to the self energy is

$$\frac{dX}{d\mathcal{S}^{(0)}} = \left\{ -\frac{1}{\beta} \sum_{\nu, |\omega_{\nu}| > \omega_x} \frac{1}{(i\hbar\omega_{\nu})^2} \right\} \left[ \mathcal{G}^{(2)} - \bar{\mathcal{G}}^{(2)} - \mathbf{s}^{-1} \left( \mathcal{S}^{(0)} - \Gamma \right) \mathbf{s}^{-1} \right] \quad (1.90)$$

Thus, if one finds a functional  $\Phi[G]$  of the Green's function alone, which produces the self energy as derivative,

$$\frac{\delta \Phi[G]}{\delta G} = \Sigma \quad (1.86)$$

we can produce a functional

$$\Phi[G] - X[G, \Sigma] \quad (1.87)$$

which produces Dyson's equation as stationary condition with respect to Green's function and self energy. The connection to the total energy is, however, no clear from this argument.

With the Laurent expansion of the non-interacting Green's function, Eq. B.17, we obtain

$$\mathcal{G}^{(2)} = \mathcal{S}^{-1} \left( \bar{h} - \mu \mathcal{S} - \left( \mathcal{S}^{(0)} - \Gamma \right) \right) \mathcal{S}^{-1} \quad (1.91)$$

The derivative of Eq. 1.89 with respect to the Green's function yields

$$\frac{dX}{d\mathcal{G}^{(2)}} = \left\{ -\frac{1}{\beta} \sum_{\nu, |\omega_\nu| > \omega_x} \frac{1}{(i\hbar\omega_\nu)^2} \right\} \mathcal{S}^{(0)} \quad (1.92)$$

The stationary condition with respect to the Laurent term of the Green's function yields

$$\begin{aligned} \frac{1}{\beta} \sum_{\nu, |\omega_\nu| > \omega_x} \left[ \frac{\beta \delta \Phi^{LW}}{\delta \mathbf{G}(i\omega_\nu)} - \frac{\beta \delta \Phi^{LW, HF}}{\delta \mathbf{G}(i\omega_\nu)} \right] \underbrace{\frac{\delta \mathbf{G}(i\omega_\nu)}{\delta \mathcal{G}^{(2)}}}_{\frac{1}{(i\hbar\omega_\nu)^2}} + \left\{ -\frac{1}{\beta} \sum_{\nu, |\omega_\nu| > \omega_x} \frac{1}{(i\hbar\omega_\nu)^2} \right\} \mathcal{S}^{(0)} = 0 \\ \mathcal{S}^{(0)} = \left\{ +\frac{1}{\beta} \sum_{\nu, |\omega_\nu| > \omega_x} \frac{1}{(i\hbar\omega_\nu)^2} \right\}^{-1} \left\{ \frac{1}{\beta} \sum_{\nu, |\omega_\nu| > \omega_x} \frac{1}{(i\hbar\omega_\nu)^2} \left[ \frac{\beta \delta \Phi^{LW}}{\delta \mathbf{G}(i\omega_\nu)} - \frac{\beta \delta \Phi^{LW, HF}}{\delta \mathbf{G}(i\omega_\nu)} \right] \right\} \quad (1.93) \end{aligned}$$

## 1.6 Dynamic-correlation correction in non-orthonormal basis and in Bloch representation

### DYNAMIC CORRELATION CORRECTION

$$\begin{aligned}
Q_{dyn,\beta}^{\hat{W}_2}[\rho] \stackrel{\text{Eq. 1.23}}{=} & \text{stat}_{\vec{k}} \text{stat}_{\vec{G}_{\vec{k}}, \Sigma_{dyn,R}} \left\{ \sum_R \left( \Phi_{\beta}^{LW}[\mathbf{G}_R, \hat{W}_{2,R}] - \Phi_{\beta}^{HF,LW}[\mathbf{G}_R, \hat{W}_{2,R}] \right) \right. \\
& - \frac{1}{N_k} \sum_k \frac{1}{\beta} \sum_{\nu, |\omega_{\nu}| < \omega_x} \text{Tr} \left\{ \ln \left[ \mathbf{1} - \left( (i\hbar\omega_{\nu} + \mu) \mathbf{S}_{\vec{k}} - \bar{\mathbf{h}}_{\vec{k}} \right)^{-1} \left( \Sigma_{dyn}(i\omega_{\nu}) - \Gamma_{\vec{k}} \right) \right] \right. \\
& \quad + \left( \Sigma_{dyn}(i\omega_{\nu}) - \Gamma_{\vec{k}} \right) \mathbf{G}_{\vec{k}}(i\omega_{\nu}) \\
& \quad \left. \left. + \left[ \mathbf{G}_{\vec{k}}(i\omega_{\nu}) - \left( (i\hbar\omega_{\nu} + \mu) \mathbf{S}_{\vec{k}} - \bar{\mathbf{h}}_{\vec{k}} \right)^{-1} \Gamma_{\vec{k}} \right] \right\} \right. \\
& - \left[ \frac{1}{\beta} \sum_{\nu, |\omega_{\nu}| > \omega_x} \frac{1}{(i\hbar\omega_{\nu})^2} \right] \frac{1}{N_k} \sum_k \text{Tr} \left\{ \left[ \mathcal{G}_{\vec{k}}^{(2)} - \mathbf{S}_{\vec{k}}^{-1} \left( \bar{\mathbf{h}}_{\vec{k}} - \mu \mathbf{S}_{\vec{k}} \right) \mathbf{S}_{\vec{k}}^{-1} \right] \mathcal{S}_{\vec{k}}^{(0)} \right. \\
& \quad \left. \left. - \frac{1}{2} \left[ \mathbf{S}_{\vec{k}}^{-1} \left( \mathcal{S}_{\vec{k}}^{(0)} - \Gamma_{\vec{k}} \right) \right]^2 \right\} \right\} \quad (1.94)
\end{aligned}$$

$\mathbf{G}_R$  is obtained from  $\mathbf{G}_{\vec{k}}$  via Eq. 1.80 and  $\Sigma_{dyn}$  is obtained from the local contributions  $\Sigma_{dyn,R}$  via Eq. 1.79. The overlap matrix  $\mathbf{S}_{\vec{k}}$  is obtained from Eq. 1.71.

$N_k$  is the number of k-points on the grid for the Brillouin-zone integration.

$$\bar{\mathbf{h}}_{a,b,\vec{k}} = \mu \mathbf{S}_{a,b,\vec{k}} + \sum_n \langle \chi_a | \psi'_n \rangle \frac{1}{\beta} \ln \left[ \frac{1 - f'_n}{f'_n} \right] \langle \psi'_n | \chi_b \rangle \quad (1.95)$$

where

$$\begin{aligned}
\left( \sum_{a,b} |\chi_a\rangle \rho_{\vec{k},a,b} \langle \chi_b| \right) \left( \sum_c |\chi_c\rangle \langle \pi_c | \psi'_{\vec{k},n} \rangle \right) &= \left( \sum_a |\chi_a\rangle \langle \pi_a | \psi'_{\vec{k},n} \rangle \right) f'_{\vec{k},n} \\
\Rightarrow \mathbf{S}_{\vec{k}} \boldsymbol{\rho}_{\vec{k}} \mathbf{S}_{\vec{k}} \vec{\psi}'_{\vec{k},n} &= \mathbf{S}_{\vec{k}} \vec{\psi}'_{\vec{k},n} f'_{\vec{k},n} \quad (1.96)
\end{aligned}$$

and  $\langle \chi_a | \chi_b \rangle = \mathbf{S}_{\vec{k}}$ .

In the high-frequency part, i.e. for  $|\omega_{\nu}| > \omega_x$ , the Green's function and the self energy are given by

$$\begin{aligned}
\mathbf{G}_{\vec{k}}(i\omega_{\nu}) &= \frac{1}{i\hbar\omega} \mathbf{S}_{\vec{k}}^{-1} + \frac{1}{(i\hbar\omega)^2} \mathcal{G}^{(2)} \\
\Sigma_{dyn,R}(i\omega_{\nu}) &= \mathcal{S}_{dyn,R}^{(0)} \quad (1.97)
\end{aligned}$$

The result is correct for  $N_k \rightarrow \infty$  and for  $\omega_x \rightarrow \infty$ .

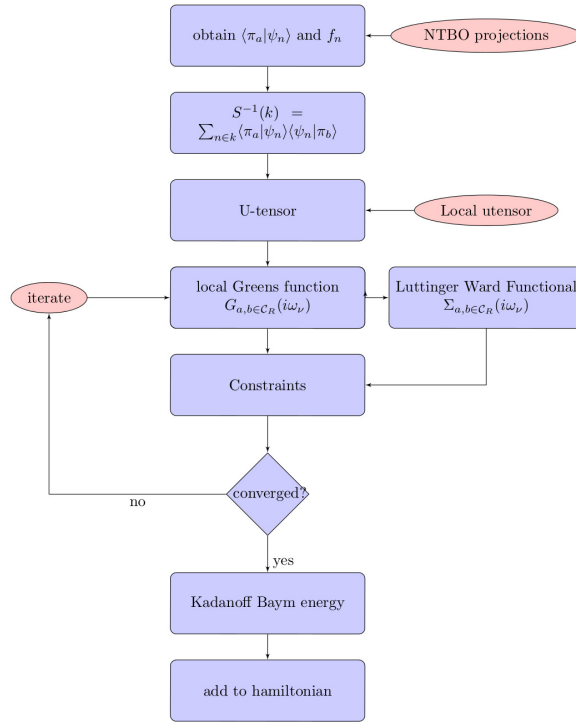
## 1.7 Description of Subroutines

### Data structures

$\langle \pi_\alpha   \psi_{n,\vec{k}} \rangle$	KSET(IKPT)%PIPSI(:NDIM,:NCHI,:NB,NSPIN)
$\rho_{\vec{k}}$	KSET(IKPT)%RHO(:NCHI,:NCHI,:NDIMD)
$S_{\vec{k}}$	KSET(IKPT)%SMAT(:NCHI,:NCHI,:NDIMD)
$S_{\vec{k}}^{-1}$	KSET(IKPT)%SINV(:NCHI,:NCHI,:NDIMD)
$\Sigma_{dyn,R}(i\omega_\nu)$	ATOMSET(IAT)%SLOC(:NLOC,:NLOC,:NDIMD,:NOMEGA)
$S_{dyn,R}^{(0)}$	ATOMSET(IAT)%SLOC1aur(:NLOC,:NLOC,:NDIMD,1)
$S_{dyn,R}^{(1)}$	ATOMSET(IAT)%SLOC1aur(:NLOC,:NLOC,:NDIMD,2)
$\bar{h}_{\vec{k}}$	KSET(IKPT)%HRHO(:NCHI,:NCHI,:NDIMD)
$\Gamma_{\vec{k}}$	KSET(IKPT)%GAMMA(:NCHI,:NCHI,:NDIMD)

#### 1.7.1 Data exchange of the Object with the outer world

#### 1.7.2 Flowchart



#### 1.7.3 DMFT\$GREEN

DMFT\$GREEN is the main subroutine of the DMFT object. It is called from the LMTO object, which also provides the projections onto the tight-binding orbitals.

We evaluate the functional  $Q_{dyn,\beta}^{W_2}[\rho]$  following the recipe provided in section III of the BPP paper[1].



```

call dmft_ini()
call dmft_collectnatorb()
call dmft_rhoofk()
call dmft_utensor()
call dmft_natorb()
call dmft_hrho()
call dmft_constraints()
do iter=1,3
  call dmft_gloc()
  call dmft_solver(etot)
  call dmft_constraints()
enddo ! end of loop over iterations to enforce constraint
call dmft_rholocal()
call dmft_staticsolver(svar)
etot=etot+svar
call dmft_detot(svar)
etot=etot+svar
call energylist$set('dmft interface',etot)
call energylist$add('local correlation',etot)
call energylist$add('total energy',etot)
call dmft_addtohpsi()

```

1. DMFT\_COLLECTNATORB: The orbital coefficients  $\langle \pi_a | \psi_n \rangle$  have been calculated by the LMTO object and they are kept in the module of `paw_waves`. The orbital coefficients are placed into `kset(ik)%pipsi(ndim,nchi,nb,nspin)`. Also the occupations are kept as `kset(ik)%f(nb,nspin)` (without spin degeneracy, i.e.  $f_n \in [0, 1]$ ). The k-point weights are kept as `kset(ik)%wkpt`. Again the spin degeneracy has been removed. However special k-points have half the weight as general k-points. **This should be checked explicitly!**
2. DMFT\_rhoofk: calculate the k-dependent density matrix, the inverse overlap matrix and the overlap matrix itself.

$$\rho_{a,b}(\vec{k}) = \sum_n \langle \pi_a | \psi_n(\vec{k}) \rangle f_n(\vec{k}) \langle \psi_n(\vec{k}) | \pi_b \rangle \quad (1.98)$$

$$\mathbf{s}_{a,b}^{-1}(\vec{k}) = \sum_n \langle \pi_a | \psi_n(\vec{k}) \rangle \langle \psi_n(\vec{k}) | \pi_b \rangle \quad (1.99)$$

$$\mathbf{s} = \left( \mathbf{s}^{-1} \right)^{-1} \quad (1.100)$$

The k-dependent density matrix is kept as `kset(ik)%rho(nchi,nchi,ndimd)`. The inverse overlap matrix is kept as `kset(ik)%sinv(nchi,nchi,ndimd)` and the overlap matrix is kept as `kset(ik)%smat(nchi,nchi,ndimd)`.

3. DMFT\_UTENSOR: (see also section 1.7.5) The local U-tensor is collected from the on-site elements `POTPAR(ISP)%TAILED%U` calculated in the `paw_LMTO` object. It is directly converted from the “tailed representation” into the basis of local orbitals.

$$U_{a,b,c,d} = \alpha \int d^4x \int d^4x' \frac{e^2 \chi_a^*(\vec{x}) \chi_b^*(\vec{x}') \chi_c(\vec{x}) \chi_d(\vec{x}')}{4\pi\epsilon_0 |\vec{r} - \vec{r}'|} \quad (1.101)$$

Only the U-tensor for equal spin-electrons is stored.  $\alpha$  is a scale factor that mimics the screening of the U-tensor. It is called LHFWEIGHT.

The scaled U-tensor is kept as `ATOMSET(iat)%U`

4. DMFT\_HRHO: (see also section 1.7.7 on p. 32.) Calculate the non-interacting Hamiltonian  $\bar{h}_{\vec{k}}$  that produces the specified (k-dependent) density matrix.

$$\bar{h}_{\vec{k},a,b} = \mu S_{\vec{k},a,b} + \sum_n \langle \chi_a | \psi'_{\vec{k},n} \rangle \frac{1}{\beta} \ln \left[ \frac{1 - f'_{\vec{k},n}}{f'_{\vec{k},n}} \right] \langle \psi'_{\vec{k},n} | \chi_b \rangle \quad (1.102)$$

$\bar{h}_{\vec{k}}$  is kept as `KSET(IKPT)%HRHO`.

5. DMFT\_NATORB: (see also section 1.7.6) Despite the name we do no more use local natural orbitals. This is described in section 1.7.6. At this point we have three different settings for the routine DMFT\_NATORB: "NATORB", "QUAMBO" and "ORTHO". QUAMBO stands for "quasi atomic molecular orbitals" [5], while it may or may not be related to it.

Construct the eigenstates  $|\phi_j\rangle$  of the single-site matrix **A**

$$\hat{A}_R |\phi_{R,i}\rangle = |\phi_{R,i}\rangle a_{R,i} \quad (1.103)$$

These states will be called local natural orbitals (This will be changed.).

The orbital coefficients  $\langle \pi_a | \phi_{R,j} \rangle$  of the local orthonormal orbitals are calculated as <sup>5</sup>

$$\sum_b \langle \pi_a | \hat{\rho}_R | \pi_b \rangle y_{b,j} = \sum_b S_{R,a,b}^{-1} y_{b,j} f'_{R,j} \quad \text{with } \vec{y}_i^\dagger \mathbf{S}_R^{-1} \vec{y}_j = \delta_{ij} \quad (1.108)$$

$$\langle \pi_a | \phi_{R,j} \rangle = \sum_b S_{R,a,b}^{-1} y_{b,j} \quad (1.109)$$

<sup>5</sup>The spirit follows from the construction of local natural orbitals along the following lines

$$\underbrace{\sum_{a,b} |\chi_a\rangle \langle \pi_a | \hat{A}_R | \pi_b \rangle \langle \chi_b |}_{\hat{\rho}_R} \underbrace{\sum_c |\chi_c\rangle \chi_{c,j}}_{|\phi_{R,j}\rangle} = \sum_a \underbrace{|\chi_a\rangle \chi_{a,j}}_{|\phi_{R,j}\rangle} f'_{R,j} \quad (1.104)$$

$$\sum_{b,c} \langle \pi_a | \hat{\rho}_R | \pi_b \rangle \langle \chi_b | \chi_c \rangle \chi_{c,j} = \chi_{a,j} f'_{R,j}$$

Now, introduce  $y_{a,j} = \sum_b \langle \chi_a | \chi_b \rangle \chi_{b,j}$

$$\sum_b \langle \pi_a | \hat{\rho}_R | \pi_b \rangle y_{b,j} = \sum_b S_{a,b}^{-1} y_{b,j} f'_{R,j} \quad (1.105)$$

which is a generalized eigenvalue problem for the vectors  $\vec{y}_j$ , which are converted into orbital coefficients by multiplication with  $\mathbf{S}^{-1}$ .

$$|\phi_{R,j}\rangle = \sum_{a,b} |\chi_a\rangle S_{a,b}^{-1} y_{b,j} \quad (1.106)$$

The local natural orbitals are orthonormal in the sense

$$\langle \phi_{R,i} | \phi_{R,j} \rangle = \sum_{a,b,c,d} y_{a,i}^* S_{a,c}^{-1} \underbrace{\langle \chi_c | \chi_d \rangle}_{(\mathbf{S}^{-1})^{-1}} S_{d,b}^{-1} y_{b,j} = \sum_{a,b} y_{a,i}^* S_{a,b}^{-1} y_{b,j} = \delta_{ij} \quad (1.107)$$

The local density matrix  $\rho_R$  and the local inverse overlap matrix  $S_R^{-1}$  are obtained as Brillouin zone integral projected onto the local site

$$\begin{aligned}\rho_R &= P_R \left[ \sum_{\vec{k}} w_{\vec{k}} \rho(\vec{k}) \right] P_R \\ S_R^{-1} &= P_R \left[ \sum_{\vec{k}} w_{\vec{k}} S^{-1}(\vec{k}) \right] P_R\end{aligned}\quad (1.110)$$

The transformation matrices are stored in the fully non-collinear data model in `atomset%natorb%piphi` and the vectors  $\vec{y}_j$  are stored as `atomset%natorb%chiphi`.

The vector  $y$  is used for the transformation to the local natural orbitals in the solver interface. Check if this is correct. should it be the vector  $x$ ?

6. DMFT\_GLOC: The local Green's function is obtained as Brillouin zone integral over the  $k$ -dependent greens function after projecting onto the correlated orbitals on the specified site.

$$G_R(i\omega_\nu) = \sum_{\vec{k}} w_{\vec{k}} \left[ (i\omega_\nu + \mu) S_{\vec{k}} - \bar{h}_{\vec{k}} + \Gamma_{\vec{k}} - \Sigma_{dyn}^{\hat{W}_2}(i\omega_\nu) \right]_R^{-1} \quad (1.111)$$

The Laurent expansion is obtained as

$$\begin{aligned}\mathcal{G}_R^{(1)}(i\omega_\nu) &= \sum_{\vec{k}} w_{\vec{k}} S_{\vec{k}}^{-1} \\ \mathcal{G}_R^{(2)}(i\omega_\nu) &= \sum_{\vec{k}} w_{\vec{k}} S_{\vec{k}}^{-1} \left( \bar{h} - \mu S_{\vec{k}} + \Gamma_{\vec{k}} - S_{dyn}^{(0)} \right) S_{\vec{k}}^{-1}\end{aligned}\quad (1.112)$$

The result and its Laurent expansion terms are stored in `atomset%gloc` and `atomset%gloc Laur`.

7. DMFT\_SOLVER: Calculates the dynamic contribution  $\Phi_{\beta}^{LW, \hat{W}_2} - \Phi_{\beta}^{HF, \hat{W}_2}$  to the Luttinger-Ward functional and the corresponding contribution  $\Sigma_{dyn}^{\hat{W}_2}$  to the self energy for the specified Green's function.

- (a) In a first loop over atoms the atoms are individually addressed and the routine `dmft_dynamicsolver` is called.
- (b) Inside `dmft_dynamicsolver` the U-tensor and the Green's function are converted into local natural orbitals and in a up-down spinor representation with all up-spin components first and the down-spin components second.
- (c) then the actual interface routine `DMFT_SOLVERIO` is called, which returns the value of the dynamic contribution to the Luttinger-Ward functional (HF subtracted), the

corresponding contribution to the self energy and the derivative of the U-tensor.

$$\begin{aligned}
\Phi_{dyn}^{LW}[\mathbf{G}(i\omega_\nu), \hat{W}_2] &= \Phi^{LW}[\mathbf{G}(i\omega_\nu), \hat{W}_2] - \frac{1}{2} \sum_{a,b,c,d} U_{a,b,d,c}^{\hat{W}_2} (\rho_{d,a} \rho_{c,b} - \rho_{c,a} \rho_{d,b}) \\
\Sigma_{dyn,a,b}^{\hat{W}_2}(i\omega_\nu) &= \frac{\beta \partial \Phi^{LW}}{\partial G_{b,a}(i\omega_\nu)} - \sum_{c,d} (U_{a,c,b,d}^{\hat{W}_2} - U_{a,c,d,b}^{\hat{W}_2}) \rho_{d,c} \\
S_{dyn,a,b}^{(0)} &= \frac{1}{\beta} \sum_{|\nu| > \nu_x} \frac{1}{(i\hbar\omega_\nu)^2} \left[ \frac{\beta \partial \Phi^{LW}}{\partial G_{b,a}(i\omega_\nu)} - \sum_{c,d} (U_{a,c,b,d}^{\hat{W}_2} - U_{a,c,d,b}^{\hat{W}_2}) \rho_{d,c} \right] \\
\frac{\delta \Phi_{\beta}^{dyn,LW}}{\delta U_{a,b,c,d}} &= \frac{\delta \Phi_{\beta}^{LW}}{\delta U_{a,b,c,d}} - \frac{1}{2} (\rho_{d,a} \rho_{c,b} - \rho_{c,a} \rho_{d,b})
\end{aligned} \tag{1.113}$$

pay attention to the reversed order of the indices  $c, d$ !

$\frac{\delta \Phi_{\beta}^{dyn,LW}}{\delta U_{a,b,c,d}}$  is the difference of the two-particle density for the interacting and the non-interacting electron gas at the same density matrix. In contrast to the corresponding term in the density functional theory, this term does not contain a kinetic energy contribution as does  $n(\vec{r}) \int_0^1 d\lambda h_\lambda(\vec{r}, \vec{r}')$ , but rather  $n(\vec{r}) [h_1(\vec{r}, \vec{r}') - h_0(\vec{r}, \vec{r}')]$

- (d) After DMFT\_SOLVERIO is finished, the self energy and the derivative with respect to the U-tensor are transformed back.

The external solver may make additional approximations to the U-tensor or to the Kadanoff-Baym functional. However, these approximations must be consistent across the dynamical-correlation correction.

8. DMFT\_MIX mixes the self energy: The self energy cannot simply be inserted into the Green's function for the next iteration, because the changes for small Matsubara frequencies would be too violent. Therefore, it needs some preconditioning.

The underlying ideas related to the mixing are described in appendix F on p. 59.

$$\begin{aligned}
\Sigma_{a,b}^{dyn}(i\omega_\nu) &\leftarrow \Sigma_{a,b}^{dyn}(i\omega_\nu) + \frac{1}{C} \left( \frac{C}{M} + (\hbar\omega)^2 \right) \left( \Sigma_{a,b}^{dyn}(i\omega_\nu) - \frac{\beta \delta \Phi_{dyn}^{LW}}{\partial G_{b,a}(i\omega_\nu)} \right) \\
S_{dyn,a,b}^{(0)} &\leftarrow S_{dyn,a,b}^{(0)} + \frac{(\hbar\omega_x)^2}{C} \left( S_{dyn,a,b}^{(0)} - \frac{\delta \Phi_{dyn}^{LW}}{\delta \mathcal{G}_{b,a}^{(2)}} \right)
\end{aligned} \tag{1.114}$$

It seems that the sign for the mixing is incorrect!! I am changing it in the code to see what happens.  $\Sigma^{dyn}$  is the self energy that has been accumulated, and that is zero in the first iteration.  $\frac{\beta \delta \Phi_{dyn}^{LW}}{\partial G_{b,a}(i\omega_\nu)}$  is the dynamic self-energy contribution produced by the solver. Note that this self energy does not contain the Hartree-Fock term.  $C$  and  $M$  are two mixing parameters, that control the behavior at small ( $M$ ) and large ( $C$ ) Matsubara frequencies.

9. DMFT\_CONSTRAINTS enforced the density matrix constraint: The new Green's function  $\mathbf{G}^{new}$  has the form

$$\mathbf{G}^{new}(i\omega_\nu) = \left[ (i\hbar\omega_\nu + \mu) \mathbf{1} - \bar{\mathbf{h}} + \mathbf{\Gamma} - \Sigma^{dyn}(i\omega_\nu) \right]^{-1} \tag{1.115}$$

where the Lagrange multiplier  $\Gamma$  needs to be adjusted until the density matrix constraint

$$\rho = \frac{1}{\beta} \sum_{\nu} e^{i\beta\hbar\omega_{\nu}0^+} \mathbf{G}^{new}(i\omega_{\nu}) \quad (1.116)$$

is fulfilled.

For this purpose, we linearize the constraint equation in the Lagrange multiplier

$$\begin{aligned} \rho &= \frac{1}{\beta} \sum_{\nu} e^{i\beta\hbar\omega_{\nu}0^+} \left[ \mathbf{G}^{new}(i\omega_{\nu}) - \mathbf{G}^{new}(i\omega_{\nu}) \delta\Gamma \mathbf{G}^{new}(i\omega_{\nu}) \right] \\ 0 &= \left[ \frac{1}{\beta} \sum_{\nu} e^{i\beta\hbar\omega_{\nu}0^+} G_{a,b}^{new}(i\omega_{\nu}) \right] - \rho_{a,b} - \sum_{c,d} \left[ \frac{1}{\beta} \sum_{\nu} \mathbf{G}_{a,c}^{new}(i\omega_{\nu}) \mathbf{G}_{d,b}^{new}(i\omega_{\nu}) \right] \delta\Gamma_{c,d} \\ \delta\Gamma_{c,d} &= \underbrace{\left[ \frac{1}{\beta} \sum_{\nu} \mathbf{G}_{a,c}^{new}(i\omega_{\nu}) \mathbf{G}_{d,b}^{new}(i\omega_{\nu}) \right]^{-1}}_{\rightsquigarrow \alpha \delta_{c,a} \delta_{d,b}} \underbrace{\left( \left[ \frac{1}{\beta} \sum_{\nu} e^{i\beta\hbar\omega_{\nu}0^+} G_{a,b}^{new}(i\omega_{\nu}) \right] - \rho_{a,b} \right)}_{\rho_{a,b}^{new}} \quad (1.117) \end{aligned}$$

In this equation we can use the model Green's function to avoid the calculation and inversion of  $\mathbf{G}^{new} \otimes \mathbf{G}^{new}$

Note, that  $\Gamma$  is, like  $\mathbf{h}'$ , a non-local Hamiltonian, that in principle connects arbitrary local orbitals with each other. In a Bloch representation,  $\Gamma$  is a  $\mathbf{k}$ -dependent matrix, which connects all local orbitals in the unit cell with each other.

Thus, we iterate the coupled equations

$$\begin{aligned} \mathbf{G}^{new}(i\omega_{\nu}) &= \left[ (i\hbar\omega_{\nu} + \mu)\mathbf{1} - \bar{\mathbf{h}} + \Gamma - \Sigma^{dyn}(i\omega_{\nu}) \right]^{-1} \\ \delta\rho &= \frac{1}{\beta} \sum_{\nu} e^{i\beta\hbar\omega_{\nu}0^+} \mathbf{G}^{new}(i\omega_{\nu}) - \rho \\ \Gamma &= \Gamma + \alpha \delta\rho \quad (1.118) \end{aligned}$$

until  $\delta\rho$  vanishes.

In practice we do not use  $\alpha$ , but precondition with a Green's function with zero one-particle level. see section 1.7.8.

With the correct Lagrange multiplier, the new Green's function is obtained.

10. DMFT\_RHLOCAL: For the Hartree-Fock and DFT double-counting terms, we need the local density matrix, which is calculated directly from the  $\mathbf{k}$ -dependent density matrix as Brillouin zone sum.
11. DMFT\_STATICSOLVER: Calculates the static contribution to the correlation energy, i.e. the screened Hartree-Fock term and the DFT double counting correction and its derivatives. These terms do not affect the density matrix constraint.

This term must be removed if the corresponding terms are treated in the PAW\_LMTO object. The double counting for the correlation must however still be included.

For the double counting in the HF approximation, only the exchange part is removed, while, here, also the correlation contribution should be taken out, because the correlation is explicitly added.

The Hartree-Fock contribution is added directly to the double counting term. Thus the self energy `atomset%sigma` is only the dynamical contribution.

12. `DMFT_DET0T`: Adds the non-local contribution to the total energy

$$-\frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \ln \left[ \mathbf{1} - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left( \hat{\Sigma}_{dyn}^{\hat{W}_2}(i\omega_{\nu}) - \mathbf{\Gamma} \right) \right. \\ \left. + \left( \hat{\Sigma}_{dyn}^{\hat{W}_2}(i\omega_{\nu}) - \mathbf{\Gamma} \right) \mathbf{G}(i\omega_{\nu}) + \left[ \mathbf{G}(i\omega_{\nu}) - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left[ \mathbf{\Gamma} \right] \right\} \quad (1.119)$$

In order to evaluate the logarithm, I diagonalize the non-hermitean matrix in the argument, and sum the absolute values of the eigenvalues. The complex logarithm is  $\ln[z] = \ln(|z|) + i\text{Arg}(z)$ . The imaginary part cancels out with the negative Matsubara frequencies.

*In a previous version I used the Dahlen trick, which is probably not quite correct. See appendix 1.7.10 on p. 36*

13. `DMFT_ADDTOHPSI`: Adds the energy derivative to `this%htbc` of the `waves_module`. There are two contributions, a k-dependent term which is restricted to the correlated orbitals and an site-dependent term that acts on all local orbitals.

- The site-specific term contains the double counting term for the DFT functional and the Hartree-Fock contribution to the self energy.

#### 1.7.4 DMFT\_INI

##### Hardwired data

- `NOMEGA` ( $N_{\omega}$ )

##### Inherited data

- `NDIM` from `waves_module`
- `NSPIN` from `waves_module`
- `NKPTL` from `waves_module`
- `NAT` from `ATOMLIST` Object
- `isp(IAT)`
- `WKPT` and `(NKPT,KMAP)` from `DYN0CC` Object
- `KBT` from `dynocc$getr8a('temp')` from `paw_occupations.f90`
- `atomset(iat)%lhweight` from `lmto_module`

**Derived data**

- NDIMD (=1 for NDIM=1,NSPIN=1; =2 for NDIM=1,NSPIN=2; =4 for NDIM=2,NSPIN=1)
- List of positive Matsubara frequencies

$$\omega_\nu = (2\nu - 1)\pi k_B T \quad \text{for } \nu = 1, \dots, N_\omega \quad (1.120)$$

- Atomset structure

atomset\%nloc    number of correlated orbits on this atom  
atomset\%ICHI1   first value of ICHI index for this atom  
atomset\%ICHI2   last value of ICHI index for this atom

- KSET structure

KSET\%WKPT    geometric k-point weight

**1.7.5 DMFT\_UTENSOR**

Obtain the on-site U-tensor for all local orbitals from subroutine `dmft_ulocal` reduces it to the entries for the correlated orbitals.

In `DMFT_ULOCAL` it obtains the local U-tensor from the `POTPAR%TAILED%U` and `SBAR` of the `LMTO` Object. This object contains the extended  $\phi$  and  $\dot{\phi}$  functions

$$|\chi_a\rangle = |\phi_a\rangle - \sum_{b; R_b=R_a} |\dot{\phi}_b\rangle \bar{S}_{a,b} \quad (1.121)$$

The U-tensor is constructed in `LMTO_MAKETAILEDPARTIALWAVES`. `LMTO_ULITTLE` constructs

$$u_{\ell,a,b,c,d} = \frac{2\ell+1}{4\pi} \int dr r^2 R_c(r) R_d(r) \left[ \int d^3 r' V_\ell(r, r') R_a(r') R_b(r') \right] \quad (1.122)$$

where the kernel  $V_\ell(r, r')$  is the one used by `RADIAL$POISSON`. **There is a input parameter that determines which  $\ell$  values are considered.**

In `LMTO_UTENSOR` the U-tensor elements are composed according to

$$U_{a,b,c,d} = \sum_{\ell} \frac{4\pi}{2\ell+1} \sum_{m=-\ell}^{\ell} u_{\ell,b,d,c,a} C_{L,L_b,L_d} C_{L,L_c,L_a} \quad (1.123)$$

The U-tensor is screened by a factor `LHFWEIGHT`. It is obtained from the `lmto_module` either as the global value `HFWEIGHT` or the individual value from `hybridsetting`. This very same factor is used for the double counting term.

Thus the U-tensor is defined as<sup>6</sup> in Eq. 1.101, i.e. as

$$U_{a,b,c,d} = \int d^4 x \int d^4 x' \frac{e^2 \chi_a^*(\vec{x}) \chi_b^*(\vec{x}') \chi_c(\vec{x}) \chi_d(\vec{x}')}{4\pi\epsilon_0 |\vec{r} - \vec{r}'|} \quad (1.124)$$

---

<sup>6</sup> $\vec{x} = (\vec{r}, \sigma)$

The U-tensor contributes only if  $\sigma_a = \sigma_c$  and if  $\sigma_b = \sigma_d$ . With this definition the Hartree and exchange energy have the form

$$\begin{aligned} E_H &= \frac{1}{2} \sum_{a,b,c,d} U_{a,b,c,d} \rho_{c,a} \rho_{d,b} \\ E_X &= \frac{1}{2} \sum_{a,b,c,d} U_{a,b,c,d} \rho_{c,b} \rho_{d,a} \end{aligned} \quad (1.125)$$

### 1.7.6 DMFT\_NATORB

Construct a set of orthonormal states for the single-site problem. This section describes first the mapping to local natural orbitals. Later we will propose a different basisset, which avoids some of the problems.

For this purpose, I redefine the unity operator for the sub Hilbert space of correlated orbitals.

$$\hat{1}_{C_R} = \sum_{\alpha,\beta} \sum_n |\chi_\alpha\rangle \langle \pi_\alpha | \psi_n \rangle \langle \psi_n | \pi_\beta \rangle \langle \chi_\beta| \quad (1.126)$$

which allows one to identify the inverse overlap matrix in this sub-Hilbert space.

$$S_{\alpha,\beta}^{-1} = \sum_{\alpha,\beta} \sum_n \langle \pi_\alpha | \psi_n \rangle \langle \psi_n | \pi_\beta \rangle \quad (1.127)$$

Note, however, that this matrix lives only on a single site! We can obtain it from the k-dependent  $S^{-1}$  by summing over k-points and cutting out the corresponding submatrix. However, it is not allowed to do the same thing for the overlap matrix itself!

We approach the problem by a formulation of natural orbitals for the local density matrix. This choice is not convenient in the end, because the resulting basisset depends on the instantaneous electronic structure. Nevertheless the point of view taken allows me to formulate the problem. Later we change it such that we use quasi atomic orbitals instead of the natural orbitals, which are less dependent on the electronic structure.

The one-particle density matrix in this sub-Hilbert space has the form

$$\hat{\rho}_{C_R} = \sum_{\alpha,\beta} \sum_n |\chi_\alpha\rangle \langle \pi_\alpha | \psi_n \rangle f_n \langle \psi_n | \pi_\beta \rangle \langle \chi_\beta| \quad (1.128)$$

The eigenvalue equation

$$\begin{aligned} \hat{\rho} |\phi_n\rangle &= |\phi_n\rangle f_n \\ \sum_{\alpha,\beta} |\chi_\alpha\rangle \langle \pi_\alpha | \hat{\rho} | \pi_\beta \rangle \langle \chi_\beta | \phi_j \rangle &= \sum_{\alpha,\beta} |\chi_\alpha\rangle \langle \pi_\alpha | \pi_\beta \rangle \langle \chi_\beta | \phi_j \rangle f_j = \sum_{\alpha,\beta} |\chi_\alpha\rangle S_{\alpha,\beta}^{-1} \langle \chi_\beta | \phi_j \rangle \bar{f}_j \end{aligned} \quad (1.129)$$

provides us with a new set of occupations  $\bar{f}$  of local natural orbitals and eigenvectors  $\langle \chi_\beta | \phi_j \rangle = V_{\beta,j}$ .

$$\rho V = S^{-1} V \bar{f} \quad \text{with} \quad V^\dagger S^{-1} V = \mathbf{1} \quad (1.130)$$

The natural orbitals are

$$|\phi_j\rangle = \sum_{\alpha,\beta \in C_R} \underbrace{|\chi_\alpha\rangle S_{\alpha,\beta}^{-1} \langle \chi_\beta|}_{\mathbf{1}_{C_R}} |\phi_j\rangle = \sum_{\alpha \in C_R} |\chi_\alpha\rangle \underbrace{(S^{-1} V)_{\alpha,j}}_{V_{\alpha,j}^{\dagger,-1}} \quad (1.131)$$



Remember that the natural orbitals are not the ones used in the calculation! The orbitals chosen are discussed in the following section.

The matrix  $\mathbf{V}$  is placed into the structure `atomset%natorb%chiphi`, together with the occupations  $\bar{f}$  which are in `atomset%natorb%f`. In addition we store  $\mathbf{S}^{-1}\mathbf{V} = \mathbf{V}^{\dagger,-1}$  in `atomset%natorb%piphi`. All data in `atomset%natorb` are stored in the fully non-collinear model.

### Orthonormal orbitals without reference to the density matrix

When approximations are done on the U-tensor, which are not invariant under a unitary transformation of the orbitals, the resulting U-tensor will be itself a functional of the natural orbitals. The approximation of a constant U-tensor in the solver will then lead to uncontrolled results.

For this reason we choose an alternative basiset instead of the local natural orbitals. The main requirement for the basiset is that it is orthonormal. Orthonormality can be obtained from any eigenvalue problem of the form

$$\mathbf{A}\mathbf{V} = \mathbf{S}^{-1}\mathbf{V}\mathbf{a} \quad (1.132)$$

where  $\mathbf{A}$  is an arbitrary hermitean matrix, the the matrix  $\mathbf{V}$  contains the eigenvectors, which obey  $\mathbf{V}^{\dagger}\mathbf{S}^{-1}\mathbf{V} = \mathbf{1}$  and the diagonal matrix  $\mathbf{a}$  contains the eigenvalues of  $\mathbf{A}$ .

This eigenvalue equation replaces Eq. 1.130.

As a second criterion, we choose orbitals that closely reflect the shape of the original tight-binding orbitals. In the current implementation the tight-binding orbitals are constructed with a clear main angular momentum character. **If the approximations made by the solver on the U-tensor are not rotationally invariant, one should use here a coordinate system that adapts to the changing local environment.**

If one chooses the matrix  $\mathbf{A}$  to be diagonal, with very different diagonal elements, then the eigenstates will be correspondingly similar to the original orbitals, and they will remain in the original order if the diagonal elements increase with index number.

Therefore, I choose  $A_{ij} = j\delta_{ij}$ . A scale factor is redundant, because it only affects the eigenvalues, but not the eigenvectors.

#### 1.7.7 DMFT\_HRHO

Calculate the non-interacting Hamiltonian  $\bar{h}$  that produces the specified (k-dependent) density matrix.

$$\bar{h}_{a,b}(\vec{k}) = \mu S_{a,b}(\vec{k}) + \sum_n \langle \chi_a | \psi'_n \rangle \frac{1}{\beta} \ln \left[ \frac{1 - f'_n}{f'_n} \right] \langle \psi'_n | \chi_b \rangle \quad (1.133)$$

where the natural orbitals  $|\psi'_n\rangle$  and the occupations  $f'_n$  differ from the original natural orbitals.

1. First, we transform onto natural orbitals

$$|\psi'_n\rangle = \sum_a |\chi_a\rangle \langle \pi_a | \psi'_n \rangle \quad (1.134)$$

These natural orbitals may be slightly different from the original natural orbitals, because they are restricted to the space of local orbitals.

$$\begin{aligned}
 \hat{\rho}|\psi'_n\rangle &= |\psi'_n\rangle f'_n \\
 \langle \chi_a | \underbrace{\sum_{b,c} |\chi_b\rangle \left( \overbrace{\sum_m \langle \pi_b | \psi_m \rangle f_m \langle \psi_m | \pi_c \rangle}^{\rho_{b,c}} \right) \langle \chi_c |}_{\hat{\rho}} \sum_d |\chi_d\rangle \langle \pi_d | \psi'_n \rangle} &= \langle \chi_a | \underbrace{\sum_d |\chi_d\rangle \langle \pi_d | \psi'_n \rangle}_{|\psi'_n\rangle} f'_n \\
 S_{a,b} \rho_{b,c} S_{c,d} \langle \pi_d | \psi'_n \rangle &= S_{a,d} \langle \pi_d | \psi'_n \rangle f'_n \\
 \boldsymbol{\rho}(\mathbf{S}\vec{x}_n) &= \mathbf{S}^{-1}(\mathbf{S}\vec{x}_n) f'_n \quad (1.135)
 \end{aligned}$$

Thus, for each k-point solve the generalized eigenvalue problem

$$\boldsymbol{\rho}(k) \mathbf{V}(\vec{k}) = \mathbf{S}^{-1}(\vec{k}) \mathbf{V}(\vec{k}) \mathbf{f}'(\vec{k})$$

where  $\mathbf{f}'(\vec{k})$  is a diagonal matrix that contains k-dependent occupations. (They differ from the occupations  $f_n$  of the natural orbitals  $|\psi_n\rangle$ .)

2. Next we evaluate the hamiltonian  $\bar{h}$  in the representation of the natural orbitals  $|\psi'_n\rangle$ .

The conversion of occupations into energies is done in the routine DMFT\_EOFF.

$$\epsilon = \mu + k_B \begin{cases} \ln\left(\frac{1-f}{f}\right) & \text{for } f_0 < f < 1 - f_0 \\ a + b(f - f_0) & \text{for } f < f_0 \\ -a + b(f - (1 - f_0)) & \text{for } f > 1 - f_0 \end{cases} \quad (1.136)$$

where  $a$  and  $b$  are determined such that the mapping is differentiable. The value of  $f_0$  is currently set to  $10^{-6}$ .

3. next we transform this Hamiltonian back into the basis of local orbitals, where the Hamiltonian has the form  $\bar{h}_{a,b} = \langle \chi_a | \hat{h} | \chi_b \rangle$ , respectively  $\hat{j} = \sum_{a,b} |\pi_a\rangle h_{a,b} \langle \pi_b|$ .

$$\begin{aligned}
 \hat{h} &= \sum_n |\psi'_n\rangle \epsilon_n \langle \psi'_n| \\
 \langle \chi_a | \hat{h} | \chi_b \rangle &= \sum_n \langle \chi_a | \psi'_n \rangle \epsilon_n \langle \psi'_n | \chi_b \rangle \\
 &= \sum_{n,c,d} \langle \chi_a | \chi_c \rangle \langle \pi_c | \psi'_n \rangle \epsilon_n \langle \psi'_n | \pi_d \rangle \langle \chi_d | \chi_b \rangle = V_{a,n} \epsilon_n V_{n,b}^\dagger \\
 \bar{h}(\vec{k}) &= \mathbf{V} \boldsymbol{\epsilon} \mathbf{V}^\dagger \quad (1.137)
 \end{aligned}$$

## 1.7.8 DMFT\_CONSTRAINTS

### Green's function

The Green's function is calculated as follows

$$\mathbf{G}(\vec{k}, i\omega_\nu) = \left[ (i\omega_\nu + \mu) \mathbf{S}(\vec{k}) - \bar{\mathbf{h}}(\vec{k}) + \boldsymbol{\Gamma}(\vec{k}) - \sum_R \boldsymbol{\Sigma}_R^{dyn}(i\omega_\nu) \right]^{-1} \quad (1.138)$$

As described in section B.2.2 on p. 46, the Laurent-expansion coefficients are

$$\begin{aligned}
 \mathcal{G}^{(1)}(\vec{k}) &= \mathbf{S}^{-1}(\vec{k}) \\
 \mathcal{G}^{(2)}(\vec{k}) &= \mathbf{S}^{-1}(\vec{k}) \left( -\mu \mathbf{S}(\vec{k}) + \bar{\mathbf{h}}(\vec{k}) - \boldsymbol{\Gamma}(\vec{k}) + \mathcal{S}_{dyn}^{(0)}(\vec{k}) \right) \mathbf{S}^{-1}(\vec{k}) \\
 \mathcal{G}^{(3)}(\vec{k}) &= \mathbf{S}^{-1}(\vec{k}) \left[ \mathcal{S}_{dyn}^{(1)}(\vec{k}) + \left( -\mu \mathbf{S}(\vec{k}) + \bar{\mathbf{h}}(\vec{k}) - \boldsymbol{\Gamma}(\vec{k}) + \mathcal{S}_{dyn}^{(0)}(\vec{k}) \right) \right. \\
 &\quad \left. \times \mathbf{S}^{-1}(\vec{k}) \left( -\mu \mathbf{S}(\vec{k}) + \bar{\mathbf{h}}(\vec{k}) - \boldsymbol{\Gamma}(\vec{k}) + \mathcal{S}_{dyn}^{(0)}(\vec{k}) \right) \right] \mathbf{S}^{-1}(\vec{k})
 \end{aligned} \tag{1.139}$$

### Green's function difference

In order to calculate the constraint, it may be advisable to directly evaluate the difference of the Green's function from the non-interacting Green's function  $\bar{\mathbf{G}}$ .

$$\begin{aligned}
 \Delta \mathcal{G}^{(1)}(\vec{k}) &= 0 \\
 \Delta \mathcal{G}^{(2)}(\vec{k}) &= \mathbf{S}^{-1}(\vec{k}) \left( \mathcal{S}_{dyn}^{(0)}(\vec{k}) - \boldsymbol{\Gamma}(\vec{k}) \right) \mathbf{S}^{-1}(\vec{k}) \\
 \Delta \mathcal{G}^{(3)}(\vec{k}) &= \mathbf{S}^{-1}(\vec{k}) \left[ \mathcal{S}_{dyn}^{(1)}(\vec{k}) + \left( -\mu \mathbf{S}(\vec{k}) + \bar{\mathbf{h}}(\vec{k}) - \boldsymbol{\Gamma}(\vec{k}) + \mathcal{S}_{dyn}^{(0)}(\vec{k}) \right) \right. \\
 &\quad \left. \times \mathbf{S}^{-1}(\vec{k}) \left( -\mu \mathbf{S}(\vec{k}) + \bar{\mathbf{h}}(\vec{k}) - \boldsymbol{\Gamma}(\vec{k}) + \mathcal{S}_{dyn}^{(0)}(\vec{k}) \right) \right. \\
 &\quad \left. - \left( -\mu \mathbf{S}(\vec{k}) + \bar{\mathbf{h}}(\vec{k}) \right) \mathbf{S}^{-1}(\vec{k}) \left( -\mu \mathbf{S}(\vec{k}) + \bar{\mathbf{h}}(\vec{k}) \right) \right] \mathbf{S}^{-1}(\vec{k})
 \end{aligned} \tag{1.140}$$

### Update $\boldsymbol{\Gamma}$

#### UPDATE $\boldsymbol{\Gamma}$

The following two equations are solved iteratively<sup>a</sup>

$$\begin{aligned}
 \delta \boldsymbol{\rho}(\vec{k}) &= \left( \frac{1}{\beta} \sum_{\nu} \underbrace{\left[ (i\omega_{\nu} + \mu) \mathbf{S}(\vec{k}) - \bar{\mathbf{h}}(\vec{k}) + \boldsymbol{\Gamma}(\vec{k}) - \sum_R \boldsymbol{\Sigma}_R^{dyn}(i\omega_{\nu}) \right]^{-1}}_{\mathbf{G}(\vec{k})} \right) \\
 &\quad - \underbrace{\left[ (i\omega_{\nu} + \mu) \mathbf{S}(\vec{k}) - \bar{\mathbf{h}}(\vec{k}) \right]^{-1}}_{\bar{\mathbf{G}}(\vec{k})} \\
 \boldsymbol{\Gamma}(\vec{k}) &= \boldsymbol{\Gamma}(\vec{k}) - \frac{4}{\beta} \mathbf{S}(\vec{k}) \delta \boldsymbol{\rho}(\vec{k}) \mathbf{S}(\vec{k})
 \end{aligned} \tag{1.141}$$

until  $\delta \boldsymbol{\rho}$  vanishes.

<sup>a</sup>Instead of subtracting the density matrix from the Matsubara sum for the Green's function, we subtract the Green's function  $\bar{\mathbf{G}}$  inside the Matsubara sum to avoid numerical problems due to the finite Matsubara sum.

This equation is a Newton scheme with the Green's function in the slope estimate replaced by the non-interacting Green's function with zero potential self energy and chemical potential.

$$\delta\rho(\Gamma) = \delta\rho(\Gamma_0) - \frac{1}{\beta} \sum_{\nu} \mathbf{G}(\Gamma - \Gamma_0) \mathbf{G} + \dots \quad (1.142)$$

Now we approximate  $\mathbf{G}$  by  $(i\hbar\omega_{\nu})^{-1} \mathbf{S}^{-1}$

$$\begin{aligned} \delta\rho(\Gamma) &= \delta\rho(\Gamma_0) - \underbrace{\frac{1}{\beta} \left[ \sum_{\nu} \frac{1}{(i\hbar\omega_{\nu})^2} \right]}_{-\beta/4} \mathbf{S}^{-1}(\Gamma - \Gamma_0) \mathbf{S}^{-1} + \dots \\ &= \delta\rho(\Gamma_0) + \frac{1}{4} \beta \mathbf{S}^{-1}(\Gamma - \Gamma_0) \mathbf{S}^{-1} + \dots \\ \Rightarrow \quad \Gamma &= \Gamma_0 + \frac{4}{\beta} \mathbf{S} \left( \underbrace{\delta\rho(\Gamma)}_{=0} - \delta\rho(\Gamma_0) \right) \mathbf{S} \end{aligned} \quad (1.143)$$

Should there be convergence problems, it may be improved by using the static Green's function  $\mathbf{G}_{\rho}$  instead.

### 1.7.9 DMFT\_SOLVER

The solver first calculates the Hartree-Fock contribution to the Luttinger-Ward functional as defined in Eq. ?? and the corresponding self energy. Because the Hartree-Fock self energy is static and because it does not affect the density matrix constraint, it is added together with the double counting correction to `atomset%denmat%h`.

Then the double counting correction is calculated. **At the moment we still subtract out only the exchange part, which is consistent with the hybrid functionals but not as double counting for correlation.**

Finally the data are prepared for the solver in `DMFT_dynamicsolver`. The `dynamicsolver` transforms all data into a spin-orbital representation of orthonormal states, which may be either natural orbitals derived from the local density matrix or which may be eigenstates of the local overlap matrix.

### 1.7.10 DMFT\_ETOT

When the loop is converged, we evaluate

$$\begin{aligned} Q_{dyn,\beta}^{\hat{W}_2}[\rho] &= \left\{ \Phi_{\beta}^{LW}[\mathbf{G}, \hat{W}_2] - \Phi_{\beta}^{HF,LW}[\mathbf{G}, \hat{W}_2] \right. \\ &\quad - \frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \ln \left[ \mathbf{1} - \bar{\mathbf{G}}(i\omega_{\nu}) \left( \Sigma_{dyn}(i\omega_{\nu}) - \Gamma \right) \right] \right. \\ &\quad \left. \left. + \left( \Sigma_{dyn}(i\omega_{\nu}) - \Gamma \right) \mathbf{G}(i\omega_{\nu}) + \left[ \mathbf{G}(i\omega_{\nu}) - \bar{\mathbf{G}}(i\omega_{\nu}) \right] \Gamma \right\} \right\} \end{aligned} \quad (1.144)$$

### Evaluation of the logarithm using a power series expansion

We use the Taylor expansion of the logarithm

$$\ln[1 - x] = - \sum_{n=1}^{\infty} \frac{1}{n} x^n \quad \text{for } |x| < 1 \quad (1.145)$$

$$\begin{aligned} Q_{dyn,\beta}^{\hat{W}_2}[\rho] &= \Phi_{\beta}^{LW}[\mathbf{G}, \hat{W}_2] - \Phi_{\beta}^{HF,LW}[\mathbf{G}, \hat{W}_2] \\ &\quad - \frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ - \sum_{n=1}^{\infty} \frac{1}{n} \left[ \bar{\mathbf{G}}(i\omega_{\nu}) (\boldsymbol{\Sigma}_{dyn}(i\omega_{\nu}) - \boldsymbol{\Gamma}) \right]^n \right. \\ &\quad \left. + (\boldsymbol{\Sigma}_{dyn}(i\omega_{\nu}) - \boldsymbol{\Gamma}) \mathbf{G}(i\omega_{\nu}) + [\mathbf{G}(i\omega_{\nu}) - \bar{\mathbf{G}}(i\omega_{\nu})] \boldsymbol{\Gamma} \right\} \\ &= \Phi_{\beta}^{LW}[\mathbf{G}, \hat{W}_2] - \Phi_{\beta}^{HF,LW}[\mathbf{G}, \hat{W}_2] \\ &\quad - \frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \boldsymbol{\Sigma}_{dyn}(i\omega_{\nu}) [\mathbf{G}(i\omega_{\nu}) - \bar{\mathbf{G}}(i\omega_{\nu})] - \sum_{n=2}^{\infty} \frac{1}{n} \left[ \bar{\mathbf{G}}(i\omega_{\nu}) (\boldsymbol{\Sigma}_{dyn}(i\omega_{\nu}) - \boldsymbol{\Gamma}) \right]^n \right\} \end{aligned} \quad (1.146)$$

### Evaluation of the logarithm using Dahlen's trick

The evaluation for the logarithm is problematic, because Green's function and self energy are not hermitean. In appendix B of their paper[6] Dahlen et al proposed a trick to solve the problem. **It is not yet clear to me if this trick is correct. See some relations for non-hermitean matrices appendix C of Bickers et al.[7]**

The trick rests on the assumption of the following identity

$$\text{Tr} \ln(\mathbf{A}\mathbf{A}^{\dagger}) \stackrel{?}{=} \text{Tr} \ln(\mathbf{A}) + \text{Tr} \ln(\mathbf{A}^{\dagger}) \quad (1.147)$$

which should be valid for arbitrary non-hermitean matrices  $\mathbf{A}$ .

Consider the singular value decomposition of  $\mathbf{A} = \mathbf{U}\boldsymbol{\Sigma}\mathbf{V}^{\dagger}$ , where  $\mathbf{U}$  and  $\mathbf{V}$  are unitary matrices and  $\boldsymbol{\Sigma}$  is a diagonal matrix with real, non-negative numbers, the singular values, on the diagonal.

$$\begin{aligned} \text{Tr} \ln(\mathbf{A}\mathbf{A}^{\dagger}) &= \text{Tr} \ln(\mathbf{U}\boldsymbol{\Sigma}\mathbf{V}^{\dagger}\mathbf{V}\boldsymbol{\Sigma}^{\dagger}\mathbf{U}^{\dagger}) = \text{Tr} \ln(\mathbf{U}\boldsymbol{\Sigma}\boldsymbol{\Sigma}^{\dagger}\mathbf{U}^{\dagger}) = \text{Tr} [\mathbf{U} \ln(\boldsymbol{\Sigma}\boldsymbol{\Sigma}^{\dagger}) \mathbf{U}^{\dagger}] \\ &= \text{Tr} [\ln(\boldsymbol{\Sigma}\boldsymbol{\Sigma}^{\dagger})] = \sum_i \ln(s_i) + \sum_i \ln(s_i) = \text{Tr} \ln[\boldsymbol{\Sigma}] + \text{Tr} \ln[\boldsymbol{\Sigma}^{\dagger}] \end{aligned}$$

$$\text{Tr} \ln(\mathbf{A}) + \text{Tr} \ln(\mathbf{A}^{\dagger}) = \text{Tr} \ln(\mathbf{U}\boldsymbol{\Sigma}\mathbf{V}^{\dagger}) + \text{Tr} \ln(\mathbf{V}\boldsymbol{\Sigma}^{\dagger}\mathbf{U}^{\dagger}) \quad (1.148)$$

Unfortunately  $\mathbf{V}^{\dagger}\mathbf{U} \neq \mathbf{1}$  so that we cannot simply convert the power series expansion into one of the singular values. **This is the problem.**

Instead of the singular value decomposition, we can also simply use the eigenvalue equation

$$\mathbf{A}\mathbf{V} = \mathbf{V}\boldsymbol{\Sigma} \quad (1.149)$$

so that

$$\begin{aligned} \mathbf{A} &= \mathbf{V}\boldsymbol{\Sigma}\mathbf{V}^{-1} \\ \mathbf{A}^{\dagger} &= \mathbf{V}^{-1,\dagger} \boldsymbol{\Sigma}^* \mathbf{V}^{\dagger} \end{aligned} \quad (1.150)$$

Unfortunately, the same problem arises, because, in general,  $\mathbf{V}\mathbf{V}^{\dagger} \neq \mathbf{1}$ .

## Chapter 2

# Tests and Numerical problems

- for insulators the Fermi-level is not well defined. It tends to fluctuate inside the band gap.
- in order to limit the number of times the solver has to be called explicitly, one may linearize the Luttinger-Ward functional

$$\Phi^{LW}[\mathbf{G}] \approx \Phi^{LW}[\mathbf{G}_0] + \frac{1}{\beta} \sum_{\nu} \sum_{a,b} \left[ \underbrace{\frac{\beta \delta \Phi^{LW}}{\delta G_{a,b}(i\omega_{\nu})}}_{\Sigma_{0,b,a}(i\omega_{\nu})} \Big|_{\mathbf{G}_0} \left( \mathbf{G}_{a,b}(i\omega_{\nu}) - \mathbf{G}_{0,a,b}(i\omega_{\nu}) \right) \right] \quad (2.1)$$

$\mathbf{G}_0$  is the last green's function for which the Luttinger ward functional and the self energy has been calculated explicitly. The approximate form can be used until convergence of the outer loop is reached. Then the Luttinger Ward functional is updated and the cycle is repeated.

## Appendix A

# Matsubara frequencies

The **Matsubara frequencies** for Fermions are

$$\omega_\nu = (2\nu + 1) \frac{\pi}{\hbar\beta}$$

A **Matsubara sum** has the general form

$$\frac{1}{\beta} \sum_{\nu \in \mathbb{Z}} g(i\omega_\nu)$$

where  $g(i\omega_\nu)$  is a analytic function.

### A.1 Evaluation using residual theorem

In order to evaluate the Matsubara sum we introduce a **Matsubara weighting functions**, that has simple poles at  $i\omega_\nu$ , so that we can exploit the residuum theorem. For Fermions, two such weighting functions are used

$$h^\pm(z) = \frac{\mp\beta\hbar}{1 + e^{\pm\beta\hbar z}} \quad (\text{A.1})$$

The poles of the Matsubara weighting functions obey

$$1 + e^{\pm\beta\hbar z} = 0 \quad (\text{A.2})$$

With the help of

$$e^{\pm\beta\hbar z} = e^{\pm\beta\hbar \text{Re}(z)} e^{\pm i\beta\hbar \text{Im}(z)} = e^{\beta\hbar \text{Re}(z)} \cos(\beta\hbar \text{Im}(z)) \pm ie^{\beta\hbar \text{Re}(z)} \sin(\beta\hbar \text{Im}(z))$$

Eq. A.2 can be rewrfitten in the form

$$\Rightarrow e^{\pm\beta\hbar \text{Re}(z)} \cos(\beta\hbar \text{Im}(z)) = -1 \quad \text{and} \quad e^{\pm\beta\hbar \text{Re}(z)} \sin(\beta\hbar \text{Im}(z)) = 0 \quad (\text{A.3})$$

The second equation is obeyed for

$$\text{Im}(z) = \frac{n\pi}{\hbar\beta} \quad \text{for arbitrary integer } n \quad (\text{A.4})$$

The first equation yields

$$\begin{aligned} -1 &= e^{\pm\beta\hbar\text{Re}(z)} \cos(\beta\hbar\text{Im}(z)) = e^{\pm\beta\hbar\text{Re}(z)} \cos(n\pi) = e^{\pm\beta\hbar\text{Re}(z)} (-1)^n \\ \Rightarrow n-1 &= 2\nu \quad \text{for arbitrary integer } \nu, \text{ and } \text{Re}(z) = 0 \\ \Rightarrow n &= 2\nu + 1 \end{aligned} \quad (\text{A.5})$$

Thus we obtain that the poles of the Matsubara weighting function lie at the Matsubara frequencies.

$$z_\nu = i(2\nu + 1) \frac{\pi}{\hbar\beta} = i\omega_\nu \quad (\text{A.6})$$

Next we need to show that the weighting function at the poles behaves like  $\frac{1}{z-z_0} + O(|z-z_0|^2)$ , where  $z_0$  is the position of the pole. For this purpose we expand the inverse

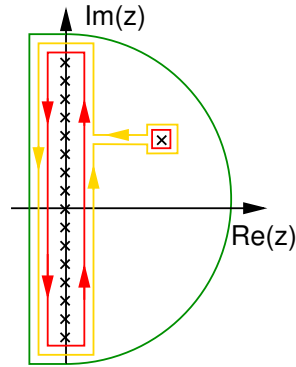
$$\frac{1}{\mp\hbar\beta} (1 + e^{\pm\hbar\beta z}) = \frac{1}{\mp\hbar\beta} \left( \underbrace{1 + e^{\pm\hbar\beta z_0}}_{=0} \pm \hbar\beta \underbrace{e^{\pm\hbar\beta z_0}}_{-1} (z - z_0) + O(|z - z_0|^2) \right) \quad (\text{A.7})$$

$$= (z - z_0) + O(|z - z_0|^2) \quad (\text{A.8})$$

With the help of the Matsubara weighting function we can evaluate the Matsubara sum as<sup>1</sup>

$$\frac{1}{\beta} \sum_{n \in \mathbb{Z}} g(i\omega_n) = \frac{1}{2\pi i\beta} \oint g(z) h^\pm(z) = -\frac{1}{\beta} \sum_{z_0 \in \text{poles of } g} \text{Res}(gh^\pm, z_0) \quad (\text{A.9})$$

The minus sign occurs because the counter clockwise integration of the contour in the half plane turns into a clock-wise integration about the pole.



The integration is closed in the half plane with  $\pm\text{Re}(z) > 0$ , because there the corresponding weighting function  $h^\pm$  decays exponentially for  $\text{Re}(z) \rightarrow \pm\infty$ .

<sup>1</sup>Cauchy's Integral formula:

$$\text{Res}(f, z_0) = \frac{1}{2\pi i} \oint_\gamma dz f(z)$$

where  $\gamma$  is a counter-clockwise contour around a pole  $z_0$  of a function  $f$ , that can be expanded into a Laurent series about  $z_0$ . The Residue is the prefactor of the term  $1/(z - z_0)$  in the Laurent expansion.



### A.1.1 Matsubara sums

From [http://en.wikipedia.org/wiki/Matsubara\\_frequency](http://en.wikipedia.org/wiki/Matsubara_frequency) we obtain the following expression for the **fermionic** summations. In these expressions the Matsubara frequencies are<sup>2</sup>

$$\omega_\nu = (2\nu + 1)\pi/(\hbar\beta) \quad (\text{A.10})$$

$$-k_B T \sum_\nu \ln[-i\hbar\omega_\nu + \epsilon] e^{i\omega_\nu 0^+} = -k_B T \ln[1 + e^{-\beta\epsilon}] \quad (\text{A.11})$$

$$k_B T \sum_\nu \frac{1}{(i\hbar\omega_\nu - \epsilon)} e^{i\omega_\nu 0^+} = (1 + e^{\beta\epsilon})^{-1} \quad (\text{A.12})$$

$$k_B T \sum_\nu \frac{1}{(i\hbar\omega_\nu - \epsilon)^n} = \frac{1}{(k_B T)^{n-1} (n-1)!} \partial_\epsilon^{n-1} \Big|_{x=\beta\epsilon} (1 + e^x)^{-1} \quad (\text{A.13})$$

$j$	1	2	3	4	5	6	7	8	9	10
$\partial_x^{j-1}(1 + e^x)^{-1}$	$\frac{1}{2}$	$-\frac{1}{4}$	0	$+\frac{1}{8}$	0	$-\frac{1}{4}$	0	$\frac{17}{16}$	0	$-\frac{31}{4}$
$\partial_x^{j-1}(1 + e^x)^{-1}$	0.5	-0.25	0	+0.125	0	-0.25	0	+1.0625	0	-7.75
$j$	11	12	13	14	15	16	17	18	19	20
$\partial_x^{j-1}(1 + e^x)^{-1}$	0	$\frac{691}{8}$	0	$-\frac{5461}{4}$	0	$\frac{929569}{32}$	0	$-\frac{3202291}{4}$	0	$+\frac{221930581}{8}$
$\partial_x^{j-1}(1 + e^x)^{-1}$	0	86.375	0	-1365.25	0	$29.0 \times 10^3$	0	$-800 \times 10^3$	0	$27.7 \times 10^6$

### A.1.2 Matsubara sums on finite grids

The finite Matsubara sums result in a Fermi function that does not reach zero or one. In particular, in the limit of  $\epsilon$  level, that lies far from the chemical potential, the finite Matsubara sum falls off to  $\frac{1}{2}$ .

<sup>2</sup>They differ from those of the source by a factor  $1/\hbar$ .

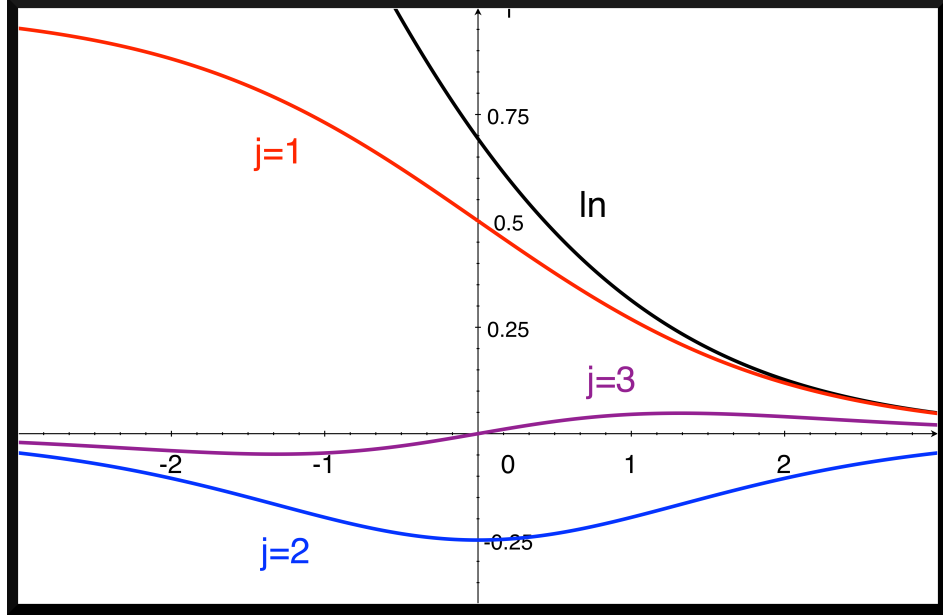


Fig. A.1: The result of Matsubara sums  $-k_B T \sum_{\nu} \ln[-i\hbar\omega_{\nu} + \epsilon]$ ,  $k_B T \sum_{\nu} \frac{1}{(i\hbar\omega_{\nu} - \epsilon)} e^{i\omega_{\nu} 0^+}$ ,  $k_B T \sum_{\nu} \frac{1}{(i\hbar\omega_{\nu} - \epsilon)^2}$  as function of energy  $\epsilon$ .

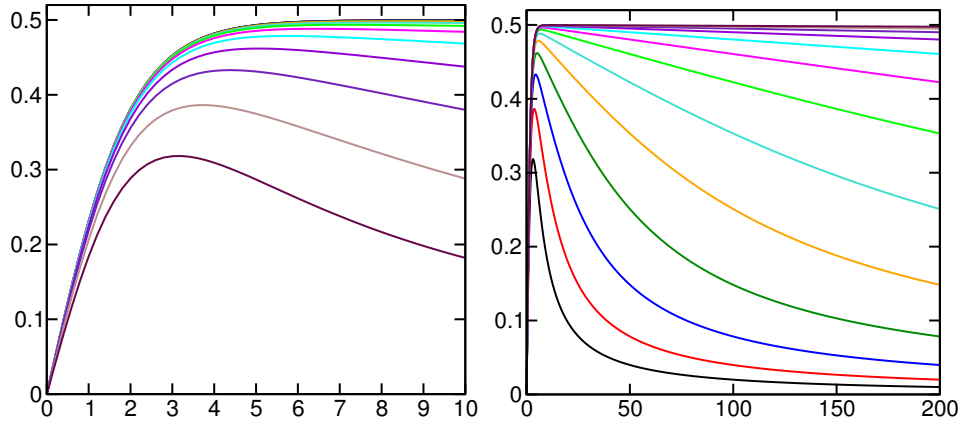


Fig. A.2: Fermi function calculated from a finite Matsubara sum. The function is shifted by  $-\frac{1}{2}$  because no regularization has been done. The sums have been performed with  $2^n$  grid points with  $n = 0, \dots, 12$  for  $k_B T = 1$ .

## Appendix B

# The Green function

We construct the **lattice Green's function** defined as

$$\begin{aligned}\hat{G}^{lat}(i\omega_\nu) &= \left[ (i\hbar\omega_\nu + \mu)\hat{1} - \hat{h} - \Delta\hat{\Sigma}(i\omega_\nu) \right]^{-1} \\ &\approx \sum_{n,n'} |\psi_n\rangle \underbrace{\left[ \langle\psi_{n'}|(i\hbar\omega_\nu + \mu)\hat{1} - \hat{h} - \Delta\hat{\Sigma}(i\omega_\nu)|\psi_n\rangle \right]_{n,n'}^{-1}}_{G_{n,n'}^{lat}} \langle\psi_{n'}| \end{aligned} \quad (\text{B.1})$$

This expression is approximate, if the set of band states is not complete. In practice, corrections are required.

In the representation of the band states, the lattice Green's function has the form

$$G_{n,n'}^{lat}(i\omega_\nu) = \left[ \langle\psi_{n'}|(i\hbar\omega_\nu + \mu)\hat{1} - \hat{h} - \Delta\hat{\Sigma}(i\omega_\nu)|\psi_n\rangle \right]_{n,n'}^{-1} \quad (\text{B.2})$$

Note, that the matrix to be inverted is not necessarily hermitean!

## B.1 Generic properties

### B.1.1 Green's function with negative Matsubara frequencies

Green's functions and self energies obey the relation

$$\mathbf{G}(-i\omega_\nu) = \mathbf{G}^\dagger(i\omega_\nu) \quad (\text{B.3})$$

This allows one to store only half of the Green's functions.

Proof:

$$G_{\alpha,\beta}(i\omega_n) = -\frac{1}{\hbar} \int_0^{\hbar\beta} d\tau e^{i\omega_n\tau} \langle \mathcal{T} \hat{c}_\alpha(\tau) \hat{c}_\beta^\dagger(0) \rangle \quad (\text{B.4})$$

$$\begin{aligned} \Rightarrow G_{\alpha,\beta}(-i\omega_n) &= -\frac{1}{\hbar} \int_0^{\hbar\beta} d\tau e^{-i\omega_n\tau} \langle \mathcal{T} \hat{c}_\alpha(\tau) \hat{c}_\beta^\dagger(0) \rangle \\ &= \left\{ -\frac{1}{\hbar} \int_0^{\hbar\beta} d\tau e^{i\omega_n\tau} \left\langle \left( e^{\beta(\hat{H}-\mu\hat{N})} \hat{c}_\alpha e^{-\beta(\hat{H}-\mu\hat{N})} \hat{c}_\beta^\dagger \right)^\dagger \right\rangle \right\}^* \\ &= \left\{ -\frac{1}{\hbar} \int_0^{\hbar\beta} d\tau e^{i\omega_n\tau} \langle \hat{c}_\beta e^{-\beta(\hat{H}-\mu\hat{N})} \hat{c}_\alpha^\dagger e^{\beta(\hat{H}-\mu\hat{N})} \rangle \right\}^* \\ &\stackrel{\text{cycl.perm}}{=} \left\{ -\frac{1}{\hbar} \int_0^{\hbar\beta} d\tau e^{i\omega_n\tau} \langle e^{\beta(\hat{H}-\mu\hat{N})} \hat{c}_\beta e^{-\beta(\hat{H}-\mu\hat{N})} \hat{c}_\alpha^\dagger \rangle \right\}^* \\ &= \left\{ -\frac{1}{\hbar} \int_0^{\hbar\beta} d\tau e^{i\omega_n\tau} \langle \mathcal{T} \hat{c}_\beta(\tau) \hat{c}_\alpha^\dagger \rangle \right\}^* \\ &= G_{\beta,\alpha}^*(i\omega_n) \end{aligned} \quad (\text{B.5})$$

### B.1.2 Density of states

For a system of independent electrons, the number of particles is related to the chemical potential via

$$N(\mu) = \int_{-\infty}^{\infty} d\epsilon f_\mu(\epsilon) D(\epsilon), \quad (\text{B.6})$$

where  $D$  is the density of states.  $f_\mu(\epsilon)$  is the Fermi function.

$$\frac{dN}{d\mu} = \int_{-\infty}^{\infty} d\epsilon \frac{\partial f_\mu(\epsilon)}{\partial \mu} D(\epsilon) = \underbrace{\int_{-\infty}^{\infty} d\epsilon \frac{\beta}{\cosh^2\left(\frac{1}{2}\beta(\epsilon - \mu)\right)} D(\epsilon)}_{\tilde{D}(\mu)} \quad (\text{B.7})$$

Thus  $dN/d\mu$  provides us with a broadened density of states.

Now we generalize the relation obtained for independent Fermions to interacting Fermions. The number of states function  $N_{a,b}(\mu)$  can be obtained from the Green's function

$$N_{a,b}(\mu) = \frac{1}{\beta} \sum_{\nu} G_{a,b}(i\omega_{\nu}, \mu) e^{i\omega_{\nu}0^+} = \frac{1}{\beta} \sum_{\nu} \left[ \left( \mathbf{G}(i\omega_{\nu}, \mu_0) \right)^{-1} + (\mu - \mu_0) \right]_{a,b}^{-1} e^{i\omega_{\nu}0^+} \quad (\text{B.8})$$

The smoothened density of states  $\tilde{D}_{a,b}(\mu)$  is then obtained as derivative with respect to  $\mu$ .

$$\begin{aligned} \tilde{D}_{a,b}(\mu) &\stackrel{\text{def}}{=} \frac{\partial N_{a,b}}{\partial \mu} \\ &= -\frac{1}{\beta} \sum_{\nu} \sum_c \left[ \left( \mathbf{G}(i\omega_{\nu}, \mu_0) \right)^{-1} + (\mu - \mu_0) \mathbf{1} \right]_{a,c}^{-1} \left[ \left( \mathbf{G}(i\omega_{\nu}, \mu_0) \right)^{-1} + (\mu - \mu_0) \mathbf{1} \right]_{c,b}^{-1} \end{aligned} \quad (\text{B.9})$$

This relation is however only valid for non-interacting electrons, because a self energy may depend itself on the chemical potential.

## B.2 Project Green's function onto local orbitals

For a Hamiltonian  $\hat{h} = \sum_{a,b} |\pi_a\rangle h_{a,b} \langle \pi_b|$  and a self energy  $\hat{\Sigma}(i\omega_\nu) = \sum_{a,b} |\pi_a\rangle \Sigma_{a,b}(i\omega_\nu) \langle \pi_b|$  acting on the local orbitals, the Green's function can be expressed as

$$\hat{G}(i\omega_\nu) = \sum_{a,b} |\chi_a\rangle \left\{ \sum_{n,n'} \langle \pi_a | \psi_n \rangle \left[ (i\hbar\omega_\nu + \mu) \delta_{n,n'} - \sum_{c,d} \langle \psi_n | \pi_c \rangle (h_{c,d} + \Sigma_{c,d}(i\omega_\nu)) \langle \pi_d | \psi_{n'} \rangle \right]_{n,n'}^{-1} \langle \psi_{n'} | \pi_b \rangle \right\} \langle \chi_b | \quad (\text{B.10})$$

We converted Eq. (B.1) into a basis of local orbitals  $|\chi_a\rangle$  using  $|\psi_n\rangle = \sum_a |\chi_a\rangle \langle \pi_a | \psi_n \rangle$ .

In order to evaluate the Green's function, we still need to refer to the band states. The Green's function can, however, also be expressed without making direct use of the band states, namely as

### GREEN'S FUNCTION IN MATRIX FORM

For a hamiltonian  $\hat{h}$  and a self energy  $\hat{\Sigma}$  that can be expressed entirely in terms of projector functions, the Green's function in the approximate form of Eq. B.1 has the form

$$\hat{G}(i\omega_\nu) = \sum_{a,b} |\chi_a\rangle G_{a,b}(i\omega_\nu) \langle \chi_b| \quad (\text{B.11})$$

with

$$\begin{aligned} & G_{a,b}(i\omega_\nu) \\ &= \left\{ \sum_{n,n'} \langle \pi_a | \psi_n \rangle \left[ (i\hbar\omega_\nu + \mu) \delta_{n,n'} - \sum_{c,d} \langle \psi_n | \pi_c \rangle (h_{c,d} + \Sigma_{c,d}(i\omega_\nu)) \langle \pi_d | \psi_{n'} \rangle \right]_{n,n'}^{-1} \langle \psi_{n'} | \pi_b \rangle \right\} \\ &= \left[ (i\hbar\omega_\nu + \mu) \mathbf{S} - \mathbf{h} - \mathbf{\Sigma}(i\omega_\nu) \right]^{-1} \end{aligned} \quad (\text{B.12})$$

where<sup>a</sup>

$$S_{a,b}^{-1} = \langle \pi_a | \psi_n \rangle \langle \psi_n | \pi_b \rangle \quad (\text{B.13})$$

<sup>a</sup>in a representation of Bloch states the inverse overlap calculation is calculated for each k-point and inverted.

**Proof:** Here, I will show that the Green's function matrix

$$\mathbf{G}_{a,b}(i\omega_\nu) = \sum_n \langle \pi_a | \psi_n \rangle \left[ (i\hbar\omega_\nu + \mu) \delta_{n,n'} - \sum_{c,d} \langle \psi_n | \pi_c \rangle (h_{c,d} + \mathbf{\Sigma}_{c,d}(i\omega_\nu)) \langle \pi_d | \psi_{n'} \rangle \right]_{n,n'}^{-1} \langle \psi_{n'} | \pi_b \rangle \quad (\text{B.14})$$

can be expressed in the form of Eq. B.12.

In order to simplify the proof, I will not write out the self energy, nor the chemical potential. This is possible because the proof works for each Matsubara frequency independently, so that

both can be absorbed in the non-interacting Hamiltonian.

$$\begin{aligned}
G_{a,b}(i\omega_\nu) &= \sum_{n,n'} \langle \pi_a | \psi_n \rangle \left( i\hbar\omega_\nu \delta_{n,n'} - \langle \psi_n | \pi \rangle \mathbf{h} \langle \pi | \psi_{n'} \rangle \right)_{n,n'}^{-1} \langle \psi_{n'} | \pi_b \rangle \\
&= \sum_{n,n'} \langle \pi_a | \psi_n \rangle \left[ \sum_{j=0}^{\infty} \frac{1}{(i\hbar\omega_\nu)^{j+1}} \left( \langle \psi | \pi \rangle \mathbf{h} \langle \pi | \psi \rangle \right)_{n,n'}^j \right] \langle \psi_{n'} | \pi_b \rangle \\
&= \sum_{n,n'} \langle \pi_a | \psi_n \rangle \left[ \sum_{j=0}^{\infty} \frac{1}{(i\hbar\omega_\nu)^{j+1}} \left( \sum_{n_2, \dots, n_j} \langle \psi_n | \pi \rangle \mathbf{h} \langle \pi | \psi_{n_2} \rangle \langle \psi_{n_2} | \pi \rangle \dots \mathbf{h} \langle \pi | \psi_{n'} \rangle \right) \right] \langle \psi_{n'} | \pi_b \rangle \\
&= \sum_{j=0}^{\infty} \frac{1}{(i\hbar\omega_\nu)^{j+1}} \left( \sum_{n_1, \dots, n_{j+1}} \underbrace{\langle \pi | \psi_{n_1} \rangle \langle \psi_{n_1} | \pi \rangle}_{\mathbf{S}^{-1}} \mathbf{h} \underbrace{\langle \pi | \psi_{n_2} \rangle \langle \psi_{n_2} | \pi \rangle}_{\mathbf{S}^{-1}} \dots \mathbf{h} \underbrace{\langle \pi | \psi_{n_{j+1}} \rangle \langle \psi_{n_{j+1}} | \pi \rangle}_{\mathbf{S}^{-1}} \right)_{a,b} \\
&= \left( \sum_{j=0}^{\infty} \frac{1}{(i\hbar\omega_\nu)^{j+1}} \left( \mathbf{S}^{-1} \mathbf{h} \right)^j \mathbf{S}^{-1} \right)_{a,b} \\
&= \left[ i\hbar\omega_\nu \mathbf{S} \left( \mathbf{1} - \frac{1}{i\hbar\omega_\nu} \mathbf{S}^{-1} \mathbf{h} \right) \right]_{a,b}^{-1} \\
&= \left[ i\hbar\omega_\nu \mathbf{S} - \mathbf{h} \right]_{a,b}^{-1}
\end{aligned} \tag{B.15}$$

q.e.d.

### B.2.1 Laurent expansion of the non-interacting Green's function

We determine the Laurent expansion terms of  $\bar{\mathbf{G}}$

$$\bar{\mathbf{G}} = \left[ i\hbar\omega_\nu \mathbf{S} - (\bar{\mathbf{h}} - \mu \mathbf{S}) \right]^{-1} = \sum_{j=1}^{\infty} \frac{1}{(i\hbar\omega_\nu)^j} \bar{\mathcal{G}}^{(j)} \tag{B.16}$$

and therefore

$$\begin{aligned}
\mathbf{1} &= \sum_{j=1}^{\infty} \frac{1}{(i\hbar\omega_\nu)^j} \bar{\mathcal{G}}^{(j)} \left[ i\hbar\omega_\nu \mathbf{S} - (\bar{\mathbf{h}} - \mu \mathbf{S}) \right] \\
&= \sum_{j=0}^{\infty} \frac{1}{(i\hbar\omega_\nu)^j} \bar{\mathcal{G}}^{(j+1)} \mathbf{S} - \sum_{j=1}^{\infty} \frac{1}{(i\hbar\omega_\nu)^j} \bar{\mathcal{G}}^{(j)} (\bar{\mathbf{h}} - \mu \mathbf{S}) \\
\bar{\mathcal{G}}^{(1)} &= \mathbf{S}^{-1} \\
\bar{\mathcal{G}}^{(j+1)} &= \bar{\mathcal{G}}^{(j)} (\bar{\mathbf{h}} - \mu \mathbf{S}) \mathbf{S}^{-1} \quad \text{for } j > 1 \\
\bar{\mathcal{G}}^j &= \mathbf{S}^{-1} \left[ (\bar{\mathbf{h}} - \mu \mathbf{S}) \mathbf{S}^{-1} \right]^{j-1}
\end{aligned} \tag{B.17}$$

The first term in the Laurent expansion of  $\mathbf{G} - \bar{\mathbf{G}}$  vanishes

$$\mathcal{G}^{(1)} - \bar{\mathcal{G}}^{(1)} = 0 \tag{B.18}$$

### B.2.2 Laurent expansion of the Green's function (new)

In order to regularize the Matsubara sums, respectively, to limit the extent of the Matsubara sum evaluated numerically, we need to represent the Green's function for high frequencies in terms of Matsubara sums.

Here, we extract the Laurent expansion of the Green's function in the form

$$\mathbf{G} = \left[ (i\hbar\omega_\nu + \mu)\mathbf{S} - \mathbf{h} - \boldsymbol{\Sigma}(i\omega_\nu) \right]^{-1} \quad (\text{B.19})$$

where the self energy has the Laurent expansion

$$\boldsymbol{\Sigma}(i\omega_\nu) = \sum_{j=0}^{\infty} \frac{1}{(i\hbar\omega_\nu)^j} \mathbf{S}^{(j)} \quad (\text{B.20})$$

Thus the numerical representation of the Green's function is

$$\mathbf{G}(i\omega_\nu) = \begin{cases} \left[ (i\hbar\omega_\nu + \mu)\mathbf{S} - \mathbf{h} - \boldsymbol{\Sigma}(i\omega_\nu) \right]^{-1} & \text{for } |\nu| \leq \nu_x \\ \left[ (i\hbar\omega_\nu + \mu)\mathbf{S} - \mathbf{h} - \sum_{j'=0}^{j'_x} \mathbf{S}^{(j')} \frac{1}{(i\hbar\omega_\nu)^{j'}} \right]^{-1} & \text{for } |\nu| > \nu_x \end{cases} \quad (\text{B.21})$$

The high-frequency representation of the Green's function shall itself be represented by a Laurent expansion

$$\mathbf{G}(i\omega_\nu) = \sum_{j=1}^{\infty} \mathbf{G}^{(j)} \frac{1}{(i\hbar\omega_\nu)^j} \quad (\text{B.22})$$

The identity

$$\sum_{j=1}^{\infty} \mathbf{G}^{(j)} \frac{1}{(i\hbar\omega_\nu)^j} = \left[ (i\hbar\omega_\nu + \mu)\mathbf{S} - \mathbf{h} - \sum_{j'=0}^{j'_x} \mathbf{S}^{(j')} \frac{1}{(i\hbar\omega_\nu)^{j'}} \right]^{-1} \quad (\text{B.23})$$

yields

$$\sum_{j=1}^{\infty} \mathbf{G}^{(j)} \frac{1}{(i\hbar\omega_\nu)^j} \left[ (i\hbar\omega_\nu + \mu)\mathbf{S} - \mathbf{h} - \sum_{j'=0}^{j'_x} \mathbf{S}^{(j')} \frac{1}{(i\hbar\omega_\nu)^{j'}} \right] = \mathbf{1} \quad (\text{B.24})$$

which we can be arranged in orders of inverse Matsubara frequencies

$$\begin{aligned} \sum_{j=0}^{\infty} \mathbf{G}^{(j+1)} \mathbf{S} \frac{1}{(i\hbar\omega_\nu)^j} + \sum_{j=1}^{\infty} \mathbf{G}^{(j)} \left[ \mu\mathbf{S} - \mathbf{h} - \mathbf{S}^{(0)} \right] \frac{1}{(i\hbar\omega_\nu)^j} \\ - \sum_{j=2}^{\infty} \sum_{n=1}^{j'_x} \mathbf{G}^{(j-n)} \mathbf{S}^{(n)} \frac{1}{(i\hbar\omega_\nu)^j} = \mathbf{1} \end{aligned} \quad (\text{B.25})$$

Thus, the expansion coefficients are

$$\begin{aligned}\mathcal{G}^{(1)} &= \mathcal{S}^{-1} \\ \mathcal{G}^{(j+1)} &= -\mathcal{G}^{(j)} \left[ \mu \mathcal{S} - \mathbf{h} - \mathcal{S}^{(0)} \right] \mathcal{S}^{-1} + \sum_{n=1}^{\min(j_x', j-1)} \mathcal{G}^{(j-n)} \mathcal{S}^{(n)} \mathcal{S}^{-1} \quad \text{for } j > 0\end{aligned} \quad (\text{B.26})$$

Let us make the higher-order terms more explicit

$$\begin{aligned}\mathcal{G}^{(1)} &= \mathcal{S}^{-1} \\ \mathcal{G}^{(2)} &= -\mathcal{G}^{(1)} \left[ \mu \mathcal{S} - \mathbf{h} - \mathcal{S}^{(0)} \right] \mathcal{S}^{-1} = -\mathcal{S}^{-1} \left[ \mu \mathcal{S} - \mathbf{h} - \mathcal{S}^{(0)} \right] \mathcal{S}^{-1} \\ \mathcal{G}^{(3)} &= -\mathcal{G}^{(2)} \left[ \mu \mathcal{S} - \mathbf{h} - \mathcal{S}^{(0)} \right] \mathcal{S}^{-1} + \mathcal{G}^{(1)} \mathcal{S}^{(1)} \mathcal{S}^{-1} \\ &= \mathcal{S}^{-1} \left[ \mu \mathcal{S} - \mathbf{h} - \mathcal{S}^{(0)} \right] \mathcal{S}^{-1} \left[ \mu \mathcal{S} - \mathbf{h} - \mathcal{S}^{(0)} \right] \mathcal{S}^{-1} + \mathcal{S}^{-1} \mathcal{S}^{(1)} \mathcal{S}^{-1}\end{aligned} \quad (\text{B.27})$$

In practice, we expand the Green's function up to  $\mathcal{G}^{(3)}$  and the self energy up to  $\mathcal{S}^{(1)}$ .

Thus the Green's function will be represented in the form

$$\mathbf{G}(i\omega_\nu) = \begin{cases} \left[ (i\hbar\omega_\nu + \mu) \mathcal{S} - \mathbf{h} - \boldsymbol{\Sigma}(i\omega_\nu) \right]^{-1} & \text{for } |\nu| \leq \nu_x \\ \sum_{j=1}^3 \mathcal{G}^{(j)} \frac{1}{(i\hbar\omega_\nu)^j} & \text{for } |\nu| > \nu_x \end{cases} \quad (\text{B.28})$$

where the Laurent-expansion coefficients of the Green's function are defined in [eqeq:Laurenttermsgreenappendix](#). The variational degrees of freedom for the Greens function is the Green's function for  $0 < \omega_\nu \leq \omega_x$  and the first two Laurent-expansion terms  $\mathcal{S}^{(0)}$  and  $\mathcal{S}^{(1)}$  of the self energy. The Green's function for negative Matsubara frequencies are obtained by [Eq. B.3](#),  $\mathbf{G}(-i\omega_\nu) = \mathbf{G}^\dagger(i\omega_\nu)$ .

### Laurent-expansion coefficients are Hermitean

The expansion coefficients for the Laurent expansion are hermitian. We exploit that  $\mathbf{G}(-i\omega_\nu) = \mathbf{G}^\dagger(i\omega_\nu)$

$$\begin{aligned}\sum_{j=1}^{\infty} \frac{1}{(-i\hbar\omega_\nu)^j} \mathcal{G}^{(j)} &= \mathbf{G}(-i\omega_\nu) = \mathbf{G}^\dagger(i\omega_\nu) = \sum_{j=1}^{\infty} \frac{1}{(-i\hbar\omega_\nu)^j} \left( \mathcal{G}^{(j)} \right)^\dagger \\ \Rightarrow \quad \mathcal{G}^{(j)} &= \left( \mathcal{G}^{(j)} \right)^\dagger\end{aligned} \quad (\text{B.29})$$

From this requirement we can derive that  $\mathbf{S}$ ,  $\mathbf{h}$ ,  $\mathcal{S}^{(0)}$  and  $\mathcal{S}^{(1)}$  are hermitian as well.

### Difference of Green's functions

The difference of two Green's function with different self energies has the form

$$\Delta \mathcal{G}^{(1)} = 0 \quad (\text{B.30})$$

$$\Delta \mathcal{G}^{(2)} = \mathcal{S}^{-1} \mathcal{S}^{(0)} \mathcal{S}^{-1} \quad (\text{B.31})$$

$$\begin{aligned}\Delta \mathcal{G}^{(3)} &= -\mathcal{S}^{-1} \left[ \mu \mathcal{S} - \mathbf{h} \right] \mathcal{S}^{-1} \mathcal{S}^{(0)} \mathcal{S}^{-1} - \mathcal{S}^{-1} \mathcal{S}^{(0)} \mathcal{S}^{-1} \left[ \mu \mathcal{S} - \mathbf{h} \right] \mathcal{S}^{-1} \\ &\quad + \mathcal{S}^{-1} \mathcal{S}^{(0)} \mathcal{S}^{-1} \mathcal{S}^{(0)} \mathcal{S}^{-1} + \mathcal{S}^{-1} \mathcal{S}^{(1)} \mathcal{S}^{-1}\end{aligned} \quad (\text{B.32})$$



### Laurent expansion for the density matrix

With this expression, we obtain the Matsubara sum required for the density matrix as

$$\begin{aligned}
\frac{1}{\beta} \sum_{\nu} \mathbf{G} e^{i\omega_{\nu} 0^+} &= \sum_{j=1}^{\infty} \mathcal{G}^{(j)} \frac{1}{(j-1)!} \partial_{\epsilon}^{j-1} \Big|_{\epsilon=0} (1 + e^{\beta\epsilon})^{-1} \\
&= \frac{1}{2} \mathcal{G}^{(1)} - \frac{1}{4} \beta \mathcal{G}^{(2)} - 0 \cdot \beta^2 \mathcal{G}^{(3)} + \underbrace{\left( \frac{1}{3!} \frac{1}{8} \beta^3 \mathcal{G}^{(4)} - \frac{1}{5!} \frac{1}{4} \beta^5 \mathcal{G}^{(6)} + \frac{1}{7!} 1.0? \beta^7 \mathcal{G}^{(8)} \right)}_{\text{from grapher}} + O(\beta^9) \\
&= \frac{1}{2} \mathcal{G}^{(1)} - \frac{1}{4} \beta \mathcal{G}^{(2)} + \underbrace{\left( \frac{1}{48} \beta^3 \mathcal{G}^{(4)} - \frac{1}{480} \beta^5 \mathcal{G}^{(6)} + \frac{1.0?}{4320} \beta^7 \mathcal{G}^{(8)} \right)}_{\text{from grapher}} + O(\beta^9)
\end{aligned}$$

### Derivative of the Luttinger Ward functional

For  $j' \in \{0, 1\}$

$$\begin{aligned}
\frac{\partial \Phi_{\beta}^{LW}}{\partial \mathcal{S}^{(j')}} &= \sum_{j=1}^3 \frac{1}{\beta} \sum_{\nu; |\nu| > \nu_x} \frac{\beta \partial \Phi_{\beta}^{LW}}{\partial \mathbf{G}(i\omega_{\nu})} \frac{\partial \mathbf{G}(i\omega_{\nu})}{\partial \mathcal{G}^{(j)}} \frac{\partial \mathcal{G}^{(j)}}{\partial \mathcal{S}^{(j')}} \\
&= \sum_{j=1}^3 \left[ \frac{1}{\beta} \sum_{\nu; |\nu| > \nu_x} \frac{\beta \partial \Phi_{\beta}^{LW}}{\partial \mathbf{G}(i\omega_{\nu})} \frac{1}{(i\hbar\omega_{\nu})^j} \right] \frac{\partial \mathcal{G}^{(j)}}{\partial \mathcal{S}^{(j')}} \\
\frac{\partial \Phi_{\beta}^{LW}}{\partial \mathcal{S}^{(0)}} &= \mathbf{S}^{-1} \underbrace{\left[ \frac{1}{\beta} \sum_{\nu; |\nu| > \nu_x} \frac{\beta \partial \Phi_{\beta}^{LW}}{\partial \mathbf{G}(i\omega_{\nu})} \frac{1}{(i\hbar\omega_{\nu})^2} \right]}_{\delta \Phi_{\beta, dyn}^{LW} / \delta \mathcal{G}^{(2)}} \mathbf{S}^{-1} \\
&\quad - \mathbf{S}^{-1} \underbrace{\left[ \frac{1}{\beta} \sum_{\nu; |\nu| > \nu_x} \frac{\beta \partial \Phi_{\beta}^{LW}}{\partial \mathbf{G}(i\omega_{\nu})} \frac{1}{(i\hbar\omega_{\nu})^3} \right]}_{\delta \Phi_{\beta, dyn}^{LW} / \delta \mathcal{G}^{(3)}} \mathbf{S}^{-1} [\mu \mathbf{S} - \mathbf{h} - \mathcal{S}^{(0)}] \mathbf{S}^{-1} \\
&\quad - \mathbf{S}^{-1} [\mu \mathbf{S} - \mathbf{h} - \mathcal{S}^{(0)}] \mathbf{S}^{-1} \underbrace{\left[ \frac{1}{\beta} \sum_{\nu; |\nu| > \nu_x} \frac{\beta \partial \Phi_{\beta}^{LW}}{\partial \mathbf{G}(i\omega_{\nu})} \frac{1}{(i\hbar\omega_{\nu})^3} \right]}_{\delta \Phi_{\beta, dyn}^{LW} / \delta \mathcal{G}^{(3)}} \mathbf{S}^{-1} \\
\frac{\partial \mathcal{G}^{(j)}}{\partial \mathcal{S}^{(1)}} &= \mathbf{S}^{-1} \underbrace{\left[ \frac{1}{\beta} \sum_{\nu; |\nu| > \nu_x} \frac{\beta \partial \Phi_{\beta}^{LW}}{\partial \mathbf{G}(i\omega_{\nu})} \frac{1}{(i\hbar\omega_{\nu})^3} \right]}_{\delta \Phi_{\beta, dyn}^{LW} / \delta \mathcal{G}^{(3)}} \mathbf{S}^{-1}
\end{aligned} \tag{B.33}$$

## Appendix C

# in progress: Green function in matrix form

In the previous section we have shown the expression for the Green's function in the form Eq. B.12, albeit under the strong condition that the hamiltonian can exactly be expressed in terms of the projector functions. The goal of this section is to perform the proof under weaker conditions.

The Green function of the non-interacting system has the form

$$G_{a,b}(i\omega_\nu) = \sum_n \langle \pi_a | \psi_n \rangle \frac{1}{i\hbar\omega_\nu + \mu - \epsilon_n} \langle \psi_n | \pi_b \rangle \quad (\text{C.1})$$

Note, that the sum over  $n$  includes the sum over  $k$ -points, and that the orbital indices may refer to orbitals from different unit cells. In practice the Green's function is calculated for each  $k$ -point individually, with orbitals only from one unit cell. The full Green's function can then be recovered from those  $k$ -projected Green's functions by a summation and using the corresponding phase factors.

The density matrix is defined by

$$\rho_{a,b} = \sum_n \langle \pi_a | \psi_n \rangle f(\epsilon_n) \langle \psi_n | \pi_b \rangle \quad (\text{C.2})$$

If (1) the Matsubara sum is complete and (2) the sum in the Green's function is performed over all states with non-zero occupations, the density matrix has the form

$$\begin{aligned} \rho_{a,b} &= \sum_n \langle \pi_a | \psi_n \rangle \left( \frac{1}{\beta} \sum_\nu e^{i\hbar\omega_\nu\beta 0^+} \frac{1}{i\hbar\omega_\nu + \mu - \epsilon_n} \right) \langle \psi_n | \pi_b \rangle \\ &= \frac{1}{\beta} \sum_\nu e^{i\hbar\omega_\nu\beta 0^+} \mathbf{G}(i\omega_\nu) \end{aligned} \quad (\text{C.3})$$

Let us rewrite the Green's function in terms of a matrix equation that does not involve a sum over all states, but with matrices on the space of local orbitals. As shown below, the result is

The Green's function matrix of the form

$$G_{a,b}(i\omega_\nu) = \sum_n^{N_x} \langle \pi_a | \psi_n \rangle \frac{1}{i\hbar\omega_\nu + \mu - \epsilon_n} \langle \psi_n | \pi_b \rangle \quad (\text{C.4})$$

can be expressed as

$$G_{a,b}(i\omega_\nu) = \sum_{n,n'}^{N_x} \langle \pi_a | \psi_n \rangle \left( (i\hbar\omega_\nu + \mu) \delta_{n,n'} - \sum_{c,d} \langle \psi_n | \pi_c \rangle h_{c,d} \langle \pi_d | \psi_{n'} \rangle - \Delta_{n,n'} \right)_{n,n'}^{-1} \langle \psi_{n'} | \pi_b \rangle \quad (\text{C.5})$$

with<sup>a</sup>

$$h_{c,d} \stackrel{\text{def}}{=} \sum_{a,b} \left( \sum_m^{N_x} \langle \pi_a | \psi_m \rangle \langle \psi_m | \pi_c \rangle \right)_{c,a}^{-1} \left( \sum_n^{N_x} \langle \pi_a | \psi_n \rangle \epsilon_n \langle \psi_n | \pi_b \rangle \right) \left( \sum_{m'}^{N_x} \langle \pi_d | \psi_{m'} \rangle \langle \psi_{m'} | \pi_b \rangle \right)_{b,d}^{-1} \quad (\text{C.6})$$

$$P_{n,n'} \stackrel{\text{def}}{=} \sum_{a,b} \langle \psi_n | \pi_a \rangle \left( \sum_m^{N_x} \langle \pi_b | \psi_m \rangle \langle \psi_m | \pi_a \rangle \right)_{a,b}^{-1} \langle \pi_b | \psi_{n'} \rangle \quad (\text{C.7})$$

$$\Delta \stackrel{\text{def}}{=} \epsilon - P\epsilon P \quad (\text{C.8})$$

All vectors and matrices with band indices are considered to be objects in a  $N_x$ -dimensional space.

It is shown that  $P$  is a projection operator in the space of band states, namely  $P^2 = P$ , while  $P \neq \mathbf{1}$ .

<sup>a</sup>To motivate the form for  $h$ , equate the denominator of both expressions for the Greens function, which gives

$$\begin{aligned} \epsilon_n \delta_{n,n'} &= \sum_{c,d} \langle \psi_n | \pi_c \rangle h_{c,d} \langle \pi_d | \psi_{n'} \rangle + \Delta_{n,n'} \\ \Rightarrow \sum_{n,n'} \langle \pi_a | \psi_n \rangle \left( \epsilon_n \delta_{n,n'} - \Delta_{n,n'} \right) \langle \psi_{n'} | \pi_b \rangle &= \sum_{n,n'} \sum_{c,d} \langle \pi_a | \psi_n \rangle \langle \psi_n | \pi_c \rangle h_{c,d} \langle \pi_d | \psi_{n'} \rangle \langle \psi_{n'} | \pi_b \rangle \\ \Leftrightarrow h_{c,d} &= \sum_{a,b} \left( \sum_m \langle \pi_a | \psi_m \rangle \langle \psi_m | \pi_c \rangle \right)_{c,a}^{-1} \left( \sum_{n,n'} \langle \pi_a | \psi_n \rangle \left( \epsilon_n \delta_{n,n'} - \Delta_{n,n'} \right) \langle \psi_{n'} | \pi_b \rangle \right) \left( \sum_{m'} \langle \pi_d | \psi_{m'} \rangle \langle \psi_{m'} | \pi_b \rangle \right)_{b,d}^{-1} \end{aligned}$$

If we set  $\Delta = 0$ , we obtain our choice for  $h$ .

The presence of  $\Delta$  in Eq. C.5 demonstrates an important limitation of Eq. B.12 given on p. 44.

### $\Delta$ and projection

For  $P$  defined through Eq. C.7, the identity

$$P \langle \psi | \pi \rangle = \langle \psi | \pi \rangle \quad (\text{C.9})$$

holds. This is shown as follows

$$\begin{aligned}
 \sum_{n'}^{N_x} P_{n,n'} \langle \psi_{n'} | \pi_c \rangle &\stackrel{\text{Eq. C.7}}{=} \sum_{n'}^{N_x} \underbrace{\sum_{a,b} \langle \psi_n | \pi_a \rangle \left( \sum_m^{N_x} \langle \pi_b | \psi_m \rangle \langle \psi_m | \pi_a \rangle \right)^{-1}}_{P_{n,n'}} \langle \pi_b | \psi_{n'} \rangle \langle \psi_{n'} | \pi_c \rangle \\
 &= \sum_{a,b} \langle \psi_n | \pi_a \rangle \underbrace{\left( \sum_m^{N_x} \langle \pi_b | \psi_m \rangle \langle \psi_m | \pi_a \rangle \right)^{-1} \left( \sum_{n'}^{N_x} \langle \pi_b | \psi_{n'} \rangle \langle \psi_{n'} | \pi_c \rangle \right)}_{\delta_{a,c}} = \langle \psi_n | \pi_c \rangle
 \end{aligned}$$

which proves Eq. C.7.

With Eq. C.7, we obtain

$$\langle \pi | \psi \rangle \Delta \langle \psi | \pi \rangle = 0 \quad (\text{C.10})$$

which is shown as follows:

$$\langle \pi | \psi \rangle \Delta \langle \psi | \pi \rangle \stackrel{\text{Eq. C.8}}{=} \langle \pi | \psi \rangle (\epsilon - P \epsilon P) \langle \psi | \pi \rangle \stackrel{\text{Eq. C.7}}{=} 0$$

### Hamilton matrix and projection

A useful relation is

$$\sum_{c,d} \langle \psi_n | \pi_c \rangle h_{c,d} \langle \pi_d | \psi_{n'} \rangle = \sum_m P_{n,m} \epsilon_m P_{m,n'} \quad (\text{C.11})$$

which is shown by insertion

$$\begin{aligned}
 &\sum_{c,d} \langle \psi_n | \pi_c \rangle h_{c,d} \langle \pi_d | \psi_{n'} \rangle \\
 &\stackrel{\text{Eq. C.6}}{=} \sum_{c,d} \langle \psi_n | \pi_c \rangle \times \\
 &\quad \overbrace{\sum_{a,b} \left( \sum_m \langle \pi_a | \psi_m \rangle \langle \psi_m | \pi_c \rangle \right)^{-1} \left( \sum_{n''} \langle \pi_a | \psi_{n''} \rangle \epsilon_{n''} \langle \psi_{n''} | \pi_b \rangle \right) \left( \sum_{m'} \langle \pi_d | \psi_{m'} \rangle \langle \psi_{m'} | \pi_b \rangle \right)^{-1}}^{h_{cd}} \langle \pi_d | \psi_{n'} \rangle \\
 &= \sum_{n''} \underbrace{\left[ \sum_{c,a} \langle \psi_n | \pi_c \rangle \left( \sum_m \langle \pi_a | \psi_m \rangle \langle \psi_m | \pi_c \rangle \right)^{-1} \langle \pi_a | \psi_{n''} \rangle \right]}_{P_{n,n''}} \epsilon_{n''} \\
 &\quad \times \underbrace{\left[ \sum_{b,d} \langle \psi_{n''} | \pi_b \rangle \left( \sum_{m'} \langle \pi_d | \psi_{m'} \rangle \langle \psi_{m'} | \pi_b \rangle \right)^{-1} \langle \pi_d | \psi_{n'} \rangle \right]}_{P_{n'',n'}} \\
 &\stackrel{\text{Eq. C.7}}{=} \left( \sum_m P_{n,m} \epsilon_m P_{m,n} \right)
 \end{aligned}$$

which proves Eq. C.11

### Proof of the identity between Eq. C.1 and Eq. C.5.

The identity between Eq. C.1 and Eq. C.5 is exact irrespective of the completeness of the band states or the local orbitals. This can be shown with the help of Eq. C.11. It is sufficient to show that the denominators are identical, namely that

$$\begin{aligned} \epsilon_n \delta_{n,n'} &= \underbrace{\left( \sum_m P_{n,m} \epsilon_m P_{m,n'} \right)}_{\sum_{c,d} \langle \psi_n | \pi_c \rangle h_{c,d} \langle \pi_d | \psi_{n'} \rangle} + \underbrace{\left( \epsilon_n \delta_{n,n'} - \sum_m P_{n,m} \epsilon_m P_{m,n'} \right)}_{\Delta_{n,n'}} \\ &\stackrel{\text{Eqs. C.11, C.8}}{=} \sum_{c,d} \langle \psi_n | \pi_c \rangle h_{c,d} \langle \pi_d | \psi_{n'} \rangle + \Delta_{n,n'} \end{aligned}$$

When this identity is inserted into Eq. C.1, we arrive directly at Eq. C.5. This completes the proof of the identity between Eq. C.1 and Eq. C.5.

### Self energy from the uncorrelated orbitals

We define the Green's function without the additional term  $\Delta$  as unperturbed (by the uncorrelated orbitals) Green's function. With the help of Eq. B.12 from p. 44 it can be expressed in matrix form without reference to the natural orbitals.

$$\begin{aligned} G_{0,a,b}^{-1}(i\omega_\nu) &\stackrel{\text{def}}{=} \sum_{n,n'} \langle \pi_a | \psi_n \rangle \left( (i\hbar\omega_\nu + \mu) \delta_{n,n'} - \sum_{c,d} \langle \psi_n | \pi_c \rangle h_{c,d} \langle \pi_d | \psi_{n'} \rangle \right)_{n,n'}^{-1} \langle \psi_{n'} | \pi_b \rangle \\ &\stackrel{\text{Eq. B.12}}{=} \left[ (i\hbar\omega_\nu + \mu) \mathbf{S} - \mathbf{h} \right]_{a,b}^{-1} \end{aligned} \quad (\text{C.12})$$

with

$$\mathbf{S}^{-1} = \langle \pi_a | \psi_n \rangle \langle \psi_n | \pi_b \rangle \quad (\text{C.13})$$

We can derive a self energy, which connects this Green function to the full Green's function by a Dyson's equation.

$$\begin{aligned} \mathbf{G}^{-1} &= \mathbf{G}_0^{-1} - \Sigma \\ \Sigma &= \left[ (i\hbar\omega_\nu + \mu) \mathbf{S} - \mathbf{h} \right] - \mathbf{G}^{-1} \\ &= \underbrace{\mathbf{P}(\epsilon - \mu)(\mathbf{1} - \mathbf{P})}_{\mathbf{H}_{AB}} \underbrace{\left( (\mathbf{1} - \mathbf{P}) \left( (i\hbar\omega_\nu + \mu) \mathbf{1} - \epsilon \right) (\mathbf{1} - \mathbf{P}) \right)^{-1}}_{\mathbf{G}_{BB}(i\omega_\nu)} \underbrace{(\mathbf{1} - \mathbf{P})(\epsilon - \mu)\mathbf{P}}_{\mathbf{H}_{BA}} \end{aligned} \quad (\text{C.14})$$

### Partitioning

The introduction of the projection operator  $\mathbf{P}$  allows one to break up the Hamiltonian into a part related to the correlated orbitals, a part outside and the hopping between correlated orbitals and the others.

Let me first review the partitioning of Green's functions

$$\begin{aligned} & -\mathbf{H}_{BA}\mathbf{G}_{AB} + (\mathbf{x}\mathbf{1}_{BB} - \mathbf{H}_{BB})\mathbf{G}_{BA} = \mathbf{0}_{BA} \\ \Rightarrow \quad \mathbf{G}_{BA} &= (\mathbf{x}\mathbf{1}_{BB} - \mathbf{H}_{BB})^{-1}\mathbf{H}_{BA}\mathbf{G}_{AB} \\ & (\mathbf{x}\mathbf{1}_{AA} - \mathbf{H}_{AA})\mathbf{G}_{AA} - \mathbf{H}_{AB}\mathbf{G}_{BA} = \mathbf{1}_{AA} \end{aligned} \quad (\text{C.15})$$

$$\begin{aligned} \xRightarrow{\text{Eq. C.15}} & (\mathbf{x}\mathbf{1}_{AA} - \mathbf{H}_{AA} - \mathbf{H}_{AB}(\mathbf{x}\mathbf{1}_{BB} - \mathbf{H}_{BB})^{-1}\mathbf{H}_{BA})\mathbf{G}_{AA} = \mathbf{1}_{AA} \\ \Rightarrow \quad \mathbf{G}_{AA} &= \left[ \mathbf{x}\mathbf{1}_{AA} - \mathbf{H}_{AA} - \underbrace{\mathbf{H}_{AB}(\mathbf{x}\mathbf{1}_{BB} - \mathbf{H}_{BB})^{-1}\mathbf{H}_{BA}}_{\Sigma(x)} \right]^{-1} \end{aligned} \quad (\text{C.16})$$

The relation between these subblocks and the projected quantities is non-trivial, which can be seen as follows:

$$\mathbf{P} \left[ (i\hbar\omega_\nu + \mu)\mathbf{1} - \epsilon \right]^{-1} \mathbf{P} = \langle \psi | \pi \rangle \mathbf{S} \mathbf{G}(i\omega_\nu) \mathbf{S} \langle \pi | \psi \rangle \quad (\text{C.17})$$

=====following is unchecked=====

The diagonal matrix of the eigenvalues of the band states can be partitioned in the form

$$\begin{aligned} \epsilon &= \mathbf{P}\epsilon\mathbf{P} + \underbrace{\epsilon - \mathbf{P}\epsilon\mathbf{P}}_{\Delta} \\ &= \mathbf{P}\epsilon\mathbf{P} + \mathbf{P}\epsilon(\mathbf{1} - \mathbf{P}) + (\mathbf{1} - \mathbf{P})\epsilon\mathbf{P} + (\mathbf{1} - \mathbf{P})\epsilon(\mathbf{1} - \mathbf{P}) \end{aligned} \quad (\text{C.18})$$

This suggests to do the following definitions

$$\begin{aligned} \mathbf{G}_{AA} &= \mathbf{P}\mathbf{G}\mathbf{P} \\ \mathbf{G}_{AB} &= \mathbf{P}\mathbf{G}(\mathbf{1} - \mathbf{P}) \\ \mathbf{1}_{AA} &= \mathbf{P}^2 \\ \mathbf{0}_{AB} &= \mathbf{P}(\mathbf{1} - \mathbf{P}) \\ \mathbf{H}_{AA} &= \mathbf{P}(\epsilon - \mu\mathbf{1})\mathbf{P} \end{aligned} \quad (\text{C.19})$$

Next we partition the Green's function

$$\begin{pmatrix} \mathbf{x}\mathbf{1}_{AA} - \mathbf{H}_{AA} & -\mathbf{H}_{AB} \\ -\mathbf{H}_{BA} & \mathbf{x}\mathbf{1}_{BB} - \mathbf{H}_{BB} \end{pmatrix} \begin{pmatrix} \mathbf{G}_{AA}(x) & \mathbf{G}_{AB}(x) \\ \mathbf{G}_{BA}(x) & \mathbf{G}_{BB}(x) \end{pmatrix} = \begin{pmatrix} \mathbf{1}_{AA} & \mathbf{0}_{AB} \\ \mathbf{0}_{BA} & \mathbf{1}_{BB} \end{pmatrix} \quad (\text{C.20})$$

Our goal is to find the Green's function  $\mathbf{G}_{AA}$  for the correlated orbitals.

This can now be written in the form

$$\begin{aligned} \mathbf{G}_{AA}(x) &= \left[ \mathbf{x} \underbrace{\mathbf{P}}_{\mathbf{1}_{AA}} - \underbrace{\mathbf{P}(\epsilon - \mu)\mathbf{P}}_{\mathbf{H}_{AA}} \right. \\ & \quad \left. - \underbrace{\mathbf{P}(\epsilon - \mu)(\mathbf{1} - \mathbf{P})}_{\mathbf{H}_{AB}} \left( (\mathbf{1} - \mathbf{P})(\mathbf{x}\mathbf{1} - \epsilon + \mu)(\mathbf{1} - \mathbf{P}) \right)^{-1} \underbrace{(\mathbf{1} - \mathbf{P})(\epsilon - \mu)\mathbf{P}}_{\mathbf{H}_{BA}} \right]^{-1} \end{aligned} \quad (\text{C.21})$$

Determine

$$\sum_n \langle \pi_a | \psi_n \rangle \langle \psi_n | \pi_b \rangle = \sum_c U_{a,c} \lambda_c U_{c,b}^\dagger \quad (\text{C.22})$$

with  $UU^\dagger = \mathbf{1}$ .

$$\begin{aligned} \left( \sum_n \langle \pi_a | \psi_n \rangle \langle \psi_n | \pi_b \rangle \right)_{a,b}^{-1} &= \sum_c U_{a,c} \frac{1}{\lambda_c} U_{c,b}^\dagger \\ \mathbf{P} &= \sum_{a,b} \langle \psi_n | \pi_a \rangle \left( \sum_m \langle \pi | \psi_m \rangle \langle \psi_m | \pi \rangle \right)_{a,b}^{-1} \langle \pi_b | \psi_n \rangle \\ &= \sum_c \left( \sum_a \langle \psi_n | \pi_a \rangle U_{a,c} \right) \frac{1}{\lambda_c} \left( U_{c,b}^\dagger \langle \pi_b | \psi_n \rangle \right) \end{aligned} \quad (\text{C.23})$$

## Appendix D

# Dividing the correlation contribution into a static and a dynamic part

We will divide the correlation contribution  $Q_{\beta}^{\hat{W}}$  in Eq. 1.7 on p. 4 into two separate terms. It seems that division is only possible if only one of the terms has a frequency-dependent self energy.

For this purpose, we express the Luttinger-Ward functional as a sum of two terms  $\Phi_{\beta}^{LW} = \Phi_{\beta}^{LW,1} + \Phi_{\beta}^{LW,2}$ , of which only the second has the frequency dependent self energy.

$$Q_{\beta}^{\hat{W}}[\rho] \stackrel{\text{Eq. 1.7}}{=} \text{stat}_{\mathbf{h}'} \text{stat}_{\mathbf{G}, \Sigma} \left\{ \Phi_{\beta}^{LW,1}[\mathbf{G}, \hat{W}_1] + \Phi_{\beta}^{LW,2}[\mathbf{G}, \hat{W}_2] \right. \\ \left. - \frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \ln \left[ \mathbf{1} - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \right\} \right. \\ \left. + \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \mathbf{G}(i\omega_{\nu}) - \left[ \mathbf{G}(i\omega_{\nu}) - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left( \mathbf{h}' - \bar{\mathbf{h}} \right) \right\} \quad (\text{D.1})$$

The stationary conditions are

$$0 = -\frac{\beta\delta\Phi^{LW,1}}{\delta\mathbf{G}(i\omega_{\nu})} - \frac{\beta\delta\Phi^{LW,2}}{\delta\mathbf{G}(i\omega_{\nu})} + \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) - \left( \mathbf{h}' - \bar{\mathbf{h}} \right) \\ \Rightarrow \quad \Sigma(i\omega_{\nu}) = \frac{\beta\delta\Phi^{LW,1}}{\delta\mathbf{G}(i\omega_{\nu})} + \frac{\beta\delta\Phi^{LW,2}}{\delta\mathbf{G}(i\omega_{\nu})} \\ 0 = \left[ \mathbf{1} - \bar{\mathbf{G}}(i\omega_{\nu}) \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \right]^{-1} \left( -\bar{\mathbf{G}}(i\omega_{\nu}) \right) + \mathbf{G}(i\omega_{\nu}) \quad (\text{D.2})$$

$$\Rightarrow \quad \mathbf{G}(i\omega_{\nu}) = \left[ \mathbf{1} - \bar{\mathbf{G}}(i\omega_{\nu}) \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \right]^{-1} \bar{\mathbf{G}}(i\omega_{\nu}) \\ 0 = -\frac{1}{\beta} \sum_{\nu} \left\{ \underbrace{\left[ \mathbf{1} - \bar{\mathbf{G}}(i\omega_{\nu}) \left( \mathbf{h}' + \Sigma(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \right]^{-1} \left( -\bar{\mathbf{G}}(i\omega_{\nu}) \right) + \mathbf{G}(i\omega_{\nu}) - \left( \mathbf{G}(i\omega_{\nu}) - \bar{\mathbf{G}}(i\omega_{\nu}) \right)}_{=0 \quad \text{Eq. D.2}} \right\} \\ \Rightarrow \quad \frac{1}{\beta} \sum_{\nu} \mathbf{G}(i\omega_{\nu}) e^{i\hbar\omega_{\nu}\beta 0^{+}} = \frac{1}{\beta} \sum_{\nu} \bar{\mathbf{G}}(i\omega_{\nu}) e^{i\hbar\omega_{\nu}\beta 0^{+}} = \rho \quad (\text{D.3})$$



Whenever the Luttinger-Ward functional can be expressed directly by the density matrix, the self energy is independent of the Matsubara frequencies. In this case the Luttinger-Ward functional depends on the Green's function only via the integral  $\rho = \frac{1}{\beta} \sum_{\nu} \mathbf{G}(i\omega_{\nu}) e^{i\hbar\omega_{\nu}\beta 0^+}$ . This is the case for the Hartree-Fock approximation and it is the case for the Density functional theory.

In the following, we assume that  $\Phi_{\beta}^{LW,1}$  has this form, i.e. that it depends on the Green's function only via the density matrix.

### Divided form

We rewrite the expression as

$$\begin{aligned} Q_{\beta}^{\hat{W}}[\rho] &= \Phi_{\beta}^{LW,1}[\bar{\mathbf{G}}, \hat{W}_1] + \text{stat}_{\mathbf{h}'_2} \text{stat}_{\mathbf{G}, \Sigma_2} \left\{ \Phi_{\beta}^{LW,2}[\mathbf{G}, \hat{W}_2] \right. \\ &\quad - \frac{1}{\beta} \sum_{\nu} \text{Tr} \left\{ \ln \left[ \mathbf{1} - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left( \mathbf{h}'_2 + \Sigma_2(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \right\} \\ &\quad \left. + \left( \mathbf{h}'_2 + \Sigma_2(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \mathbf{G}(i\omega_{\nu}) - \left[ \mathbf{G}(i\omega_{\nu}) - \left( i\hbar\omega_{\nu} + \mu \right) \mathbf{1} - \bar{\mathbf{h}} \right]^{-1} \left( \mathbf{h}'_2 - \bar{\mathbf{h}} \right) \right\} \end{aligned} \quad (\text{D.4})$$

The stationary conditions are

$$\begin{aligned} \Sigma_2(i\omega_{\nu}) &= \frac{\beta \delta \Phi^{LW,2}}{\delta \mathbf{G}(i\omega_{\nu})} \\ \mathbf{G}(i\omega_{\nu}) &= \left[ \mathbf{1} - \bar{\mathbf{G}}(i\omega_{\nu}) \left( \mathbf{h}'_2 + \Sigma_2(i\omega_{\nu}) - \bar{\mathbf{h}} \right) \right]^{-1} \bar{\mathbf{G}}(i\omega_{\nu}) \end{aligned} \quad (\text{D.5})$$

We see that

$$\mathbf{h}' = \mathbf{h}'_2 + \frac{\beta \delta \Phi^{LW,1}}{\delta \mathbf{G}(i\omega_{\nu})} \quad (\text{D.6})$$

leads to the same Green's function as in the former case. That is, both forms, i.e. Eq. D.1 and Eq. D.4, have the same stationary state. At the stationary state both forms have furthermore the same value.

### No Matsubara sum with a stationary self energy

Here, I show that  $Q_{\beta}^{\hat{W}}[\rho] = \Phi_{\beta}^{LW}[\bar{\mathbf{G}}, \hat{W}]$ , if the Luttinger Ward functional only depends on the Greens function via the density matrix.

In that case, the self energy is frequency independent. All stationary conditions can be fulfilled by choosing

$$\begin{aligned} \mathbf{h}' + \Sigma &= \bar{\mathbf{h}} \\ \mathbf{G}(i\omega_{\nu}) &= \bar{\mathbf{G}}(i\omega_{\nu}) \end{aligned} \quad (\text{D.7})$$

Thus we may decompose the correlation contribution into several terms of which only one has a frequency dependent self energy

## Appendix E

# Hartree-Fock energy

. The Hartree-Fock term  $Q_{X,\beta}^{\hat{W}}$  is equal to  $Q_{\beta}^{\hat{W}}$  when only the first-order term of the Luttinger-Ward functional in the interaction is considered. It is obtained as (BPP-Eq.43)

$$Q_{X,\beta}^{\hat{W}}[\rho] = \frac{1}{2} \sum_{a,b,c,d} U_{a,b,d,c} [\rho_{d,a} \rho_{c,b} - \rho_{c,a} \rho_{d,b}] \quad (\text{E.1})$$

The interpretation of this term is subtle because it is formulated in non-orthonormal orbitals.

- We consider the expansion of Kohn-Sham orbitals in local orbitals

$$\begin{aligned} Q_{X,\beta}^{\hat{W}}[\rho] &= \frac{1}{2} \sum_{m,n} f_m f_n \int d^3 r \int d^3 r' \frac{e^2 (\phi_m^*(\vec{r}) \phi_n^*(\vec{r}') \phi_n(\vec{r}') \phi_m(\vec{r}) - \phi_m^*(\vec{r}) \phi_n^*(\vec{r}') \phi_m(\vec{r}') \phi_n(\vec{r}))}{4\pi\epsilon_0 |\vec{r} - \vec{r}'|} \\ &= \frac{1}{2} \sum_{a,b,c,d} \underbrace{\left( \sum_m \langle \pi_a | \phi_m \rangle f_m \langle \phi_m | \pi_b \rangle \right)}_{\rho_{a,b}} \underbrace{\left( \sum_n \langle \pi_c | \phi_n \rangle f_n \langle \phi_n | \pi_d \rangle \right)}_{\rho_{c,d}} \\ &\quad \times \int d^3 r \int d^3 r' \frac{e^2 (\chi_b^*(\vec{r}) \chi_d^*(\vec{r}') \chi_c(\vec{r}') \chi_a(\vec{r}) - \chi_b^*(\vec{r}) \chi_d^*(\vec{r}') \chi_a(\vec{r}') \chi_c(\vec{r}))}{4\pi\epsilon_0 |\vec{r} - \vec{r}'|} \end{aligned} \quad (\text{E.2})$$

$$= \frac{1}{2} \sum_{a,b,c,d} \rho_{a,b} \rho_{c,d} (U_{b,d,a,c} - U_{b,d,c,a}) \quad (\text{E.3})$$

$$= \frac{1}{2} \sum_{a,b,c,d} U_{a,b,d,c} (\rho_{d,a} \rho_{c,b} - \rho_{c,a} \rho_{d,b}) \quad (\text{E.4})$$

- Consider expectation value of a Slater determinant expressed in terms of non-orthonormal local orbitals. The Slater determinant has the form

$$\Psi(\vec{x}_1, \dots, \vec{x}_N) = C \det |\mathbf{M}| = C \sum_{i_1, \dots, i_N=1}^N \epsilon_{i_1, \dots, i_N} \chi_{i_1}(\vec{x}_1) \cdots \chi_{i_N}(\vec{x}_N) \quad (\text{E.5})$$

with  $M_{i,j} = \chi_i(\vec{x}_j)$  and the normalization constant  $C$ .  $\epsilon_{i_1, \dots, i_N}$  is the fully antisymmetric tensor defined by  $\epsilon_{1,2,\dots,N} = 1$ , by the fact that it changes sign under permutation of two indices, and that it vanishes whenever two indices are identical.

This argument needs to be completed. Probably we obtain the same result as in the first case. This needs to be shown by performing a transformation onto orthogonal one-particle states, that span the same Hilbert space.

## Appendix F

# Mixing of the self energy

### F.1 Structure of the self-consistency

The self-consistent scheme has the following structure

$$\Sigma_{dyn}(+) = \Sigma_{dyn}(0) + \alpha \left[ \frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta \mathbf{G}} \Big|_{\mathbf{G}[\Sigma_{dyn}(0)]} - \Sigma_{dyn}(0) \right] \quad (\text{F.1})$$

and we search for the fixed point of this dynamics.  $\mathbf{G}[\Sigma_{dyn}]$  is given by Dyson's equation, that is by the stationary condition of the dynamic-correlation correction with respect to the self energy.

This problem is equivalent to the search of a zero of  $\frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta \mathbf{G}} - \Sigma_{dyn}$ . More precisely, we search  $\delta \Sigma_{dyn}$  so that

$$\frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta \mathbf{G}} \Big|_{\mathbf{G}[\Sigma_{dyn} + \delta \Sigma_{dyn}]} - \Sigma_{dyn} - \delta \Sigma_{dyn} = 0 \quad (\text{F.2})$$

This yields to first order in  $\delta \Sigma_{dyn}$

$$\left[ \frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta \mathbf{G}} \Big|_{\mathbf{G}[\Sigma_{dyn}]} - \Sigma_{dyn} \right] + \frac{\beta \delta^2 \Phi_{\beta}^{LW,dyn}}{\delta \mathbf{G} \delta \mathbf{G}'} \Big|_{\mathbf{G}[\Sigma_{dyn}]} \frac{\delta \mathbf{G}'}{\delta \Sigma_{dyn}} \delta \Sigma_{dyn} - \delta \Sigma_{dyn} = 0 \quad (\text{F.3})$$

I write out the indices

$$\left[ \frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta G_{b,a}} \Big|_{\mathbf{G}[\Sigma_{dyn}]} - \Sigma_{dyn,a,b} \right] + \sum_{e,f} \left\{ \left[ \sum_{c,d} \frac{\beta \delta^2 \Phi_{\beta}^{LW,dyn}}{\delta G_{b,a} \delta G_{d,c}} \Big|_{\mathbf{G}[\Sigma_{dyn}]} \frac{\delta G_{d,c}}{\delta \Sigma_{dyn,e,f}} \right] - \delta_{a,e} \delta_{b,f} \right\} \delta \Sigma_{dyn,e,f} = 0 \quad (\text{F.4})$$

We can now evaluate the derivative of the Green's function with respect to the self energy.

$$\begin{aligned} \mathbf{G} &= \left[ (i\hbar\omega_{\nu} + \mu)\mathbf{S} - \bar{\mathbf{h}} - (\Sigma_{dyn} - \mathbf{\Gamma}) \right]^{-1} \\ \Rightarrow \quad \delta \mathbf{G} &= \mathbf{G} (\delta \Sigma_{dyn} - \delta \mathbf{\Gamma}) \mathbf{G} \\ \Rightarrow \quad \delta G_{d,c} &= G_{d,e} (\delta \Sigma_{dyn,e,f} - \delta \Gamma_{e,f}) G_{f,c} \end{aligned} \quad (\text{F.5})$$

and obtain

$$\left[ \frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta G_{b,a}} \right]_{\mathbf{G}[\Sigma_{dyn}]} - \Sigma_{dyn,a,b} + \sum_{e,f} \left\{ \left[ \sum_{c,d} G_{c,e} \frac{\beta \delta^2 \Phi_{\beta}^{LW,dyn}}{\delta G_{b,a} \delta G_{c,d}} \right]_{\mathbf{G}[\Sigma_{dyn}]} G_{f,d} - \delta_{a,e} \delta_{b,f} \right\} \delta \Sigma_{dyn,e,f} = 0 \quad (\text{F.6})$$

In order to devise an efficient preconditioning scheme, we need a model Green's function and a model dynamical Luttinger-Ward functional. A model Green's function is developed below in section F.3 on p. 61.

$$G_{a,b}(i\omega_{\nu}) G_{c,d}(i\omega_{\nu}) \approx -\mathbf{S}_{a,b}^{-1} \mathbf{S}_{c,d}^{-1} \frac{1}{A + B(\hbar\omega_{\nu})^2} \quad (\text{F.7})$$

Without a good approximation for the Luttinger Ward functional, we make the following simple approximation

$$\sum_{c,d} G_{c,e} \frac{\beta \delta^2 \Phi_{\beta}^{LW,dyn}}{\delta G_{b,a} \delta G_{c,d}} \Big|_{\mathbf{G}[\Sigma_{dyn}]} G_{f,d} \approx -\delta_{a,e} \delta_{b,f} \frac{1}{A + B(\hbar\omega_{\nu})^2} \quad (\text{F.8})$$

With this approximation, I obtain the prediction for  $\delta \Sigma_{dyn}$  as

$$\begin{aligned} \left[ \frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta G_{b,a}} \right]_{\mathbf{G}[\Sigma_{dyn}]} - \Sigma_{dyn,a,b} - \sum_{e,f} \left\{ \delta_{a,e} \delta_{b,f} \frac{1}{A + B(\hbar\omega_{\nu})^2} + \delta_{a,e} \delta_{b,f} \right\} \delta \Sigma_{dyn,e,f} &= 0 \\ \left[ \frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta G_{b,a}} \right]_{\mathbf{G}[\Sigma_{dyn}]} - \Sigma_{dyn,a,b} + \left\{ 1 + \frac{1}{A + B(\hbar\omega_{\nu})^2} \right\} \delta \Sigma_{dyn,a,b} &= 0 \end{aligned} \quad (\text{F.9})$$

so that

$$\delta \Sigma_{dyn,a,b} \approx \frac{1}{1 + \frac{1}{A + B(\hbar\omega_{\nu})^2}} \left[ \frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta G_{b,a}} \right]_{\mathbf{G}[\Sigma_{dyn}]} - \Sigma_{dyn,a,b} \quad (\text{F.10})$$

Thus, in the high-frequency limit, it seems that a mixing factor equal to one seems to be ideal, while the low frequency limit should have an independent mixing.

A more intuitive set of parameters that  $A, B$  is to specify the mixing  $\alpha$  for  $\omega = 0$  and the Matsubara frequency  $\omega_h$  for which the mixing is halfway between  $\alpha$  and one.

$$\Sigma_{dyn}(+) = \Sigma_{dyn}(0) + \frac{1}{1 + \frac{1-\alpha}{\alpha} + \left(\frac{\omega_{\nu}}{\omega_h}\right)^2} \left[ \frac{\beta \delta \Phi_{\beta}^{LW,dyn}}{\delta G_{b,a}} \right]_{\mathbf{G}[\Sigma_{dyn}(0)]} - \Sigma_{dyn,a,b}(0) \quad (\text{F.11})$$

The Laurent expansions are mixed with a mixing factor equal to unity consistent with the high-frequency mixing.

## F.2 Euler-Lagrange equations

In order to arrive at a reasonable mixing scheme, we explore a dynamics from a Lagrangian formalism. It is however important to note that this dynamics is useless, because there is no minimum principle, so that applying a friction will not lead to a stationary state.

$$m_{\Sigma}(\omega_{\nu})\ddot{\Sigma}_{a,b}(i\omega_{\nu}) = -\beta \frac{\partial E}{\partial \Sigma_{a,b}^{\dagger}(i\omega_{\nu})}$$

The equation for the self energy is a little complicated. Therefore we obtain it directly from the action principle.

$$\begin{aligned} \mathcal{L} &= \frac{1}{\beta} \sum_{\nu} m_{\Sigma} \sum_{a,b} \left\{ \dot{\Sigma}_{a,b}(i\omega_{\nu}) \dot{\Sigma}_{b,a}^{\dagger}(i\omega_{\nu}) \right\} - \Phi(\mathbf{G}) \\ 0 &= \frac{1}{\beta} \sum_{\nu} m_{\Sigma} \sum_{a,b} \left\{ \dot{\Sigma}_{a,b}^{\dagger}(i\omega_{\nu}) \delta \dot{\Sigma}_{b,a}(i\omega_{\nu}) + \dot{\Sigma}_{a,b}(i\omega_{\nu}) \delta \dot{\Sigma}_{b,a}^{\dagger}(i\omega_{\nu}) \right\} \\ &\quad - \frac{1}{\beta} \sum_{\nu} \sum_{a,b,c,d} \left[ \frac{\beta \partial \Phi}{\partial G_{a,b}(i\omega_{\nu})} \frac{\partial G_{a,b}(i\omega_{\nu})}{\partial \Sigma_{c,d}(i\omega_{\nu})} \delta \Sigma_{c,d}(i\omega_{\nu}) + \frac{\beta \partial \Phi}{\partial G_{a,b}^{\dagger}(i\omega_{\nu})} \frac{\partial G_{a,b}^{\dagger}(i\omega_{\nu})}{\partial \Sigma_{c,d}^{\dagger}(i\omega_{\nu})} \delta \Sigma_{c,d}^{\dagger}(i\omega_{\nu}) \right] \\ &= \frac{1}{\beta} \sum_{\nu} m_{\Sigma} \sum_{a,b} \left\{ -\ddot{\Sigma}_{a,b}^{\dagger}(i\omega_{\nu}) \delta \Sigma_{b,a}(i\omega_{\nu}) - \ddot{\Sigma}_{a,b}(i\omega_{\nu}) \delta \Sigma_{b,a}^{\dagger}(i\omega_{\nu}) \right\} \\ &\quad - \frac{1}{\beta} \sum_{\nu} \sum_{a,b,c,d} \frac{\beta \partial \Phi}{\partial G_{a,b}(i\omega_{\nu})} G_{a,c}(i\omega_{\nu}) \delta \Sigma_{c,d}(i\omega_{\nu}) G_{d,b}(i\omega_{\nu}) \\ &\quad - \frac{1}{\beta} \sum_{\nu} \sum_{a,b,c,d} \frac{\beta \partial \Phi}{\partial G_{a,b}^{\dagger}(i\omega_{\nu})} G_{a,c}^{\dagger}(i\omega_{\nu}) \delta \Sigma_{c,d}^{\dagger}(i\omega_{\nu}) G_{d,b}^{\dagger}(i\omega_{\nu}) + \text{boundary terms at } t_1, t_2 \\ &= -\frac{1}{\beta} \sum_{\nu} \left[ \sum_{a,b} m_{\Sigma} \ddot{\Sigma}_{a,b}(i\omega_{\nu}) \delta \Sigma_{b,a}^{\dagger}(i\omega_{\nu}) + \sum_{c,d} \left( \sum_{a,b} G_{d,b}^{\dagger}(i\omega_{\nu}) \frac{\beta \partial \Phi}{\partial G_{a,b}^{\dagger}(i\omega_{\nu})} G_{a,c}^{\dagger}(i\omega_{\nu}) \right) \delta \Sigma_{c,d}^{\dagger}(i\omega_{\nu}) \right] + \text{c.c.} \\ &= -\frac{1}{\beta} \sum_{\nu} \sum_{a,b} \left[ m_{\Sigma} \ddot{\Sigma}_{a,b}(i\omega_{\nu}) + \sum_{c,d} G_{a,c}^{\dagger}(i\omega_{\nu}) \frac{\beta \partial \Phi}{\partial G_{d,c}^{\dagger}(i\omega_{\nu})} G_{d,b}^{\dagger}(i\omega_{\nu}) \right] \delta \Sigma_{b,a}^{\dagger}(i\omega_{\nu}) + \text{c.c.} \end{aligned}$$

Thus the equation of motion has the form

$$m_{\Sigma} \ddot{\Sigma}_{a,b}(i\omega_{\nu}) = - \sum_{c,d} G_{a,c}^{\dagger}(i\omega_{\nu}) \frac{\beta \partial \Phi}{\partial G_{d,c}^{\dagger}(i\omega_{\nu})} G_{d,b}^{\dagger}(i\omega_{\nu})$$

With  $G(i\omega) = G^{\dagger}(-i\omega)$  and  $\Sigma(i\omega) = \Sigma^{\dagger}(-i\omega)$ , I obtain

$$m_{\Sigma} \ddot{\Sigma}_{a,b}^{\dagger}(i\omega_{\nu}) = - \sum_{c,d} G_{a,c}(i\omega_{\nu}) \frac{\beta \partial \Phi}{\partial G_{d,c}(i\omega_{\nu})} G_{d,b}(i\omega_{\nu})$$

### F.3 Model Green's function

In order to analyze the equations it is often useful to have a model Green's function that has an analytical form but reflects the qualitative properties of the Green's function.

We start with the expression for a non-interacting Green's function in terms of Matsubara frequencies

$$G_{a,b}(i\omega_{\nu}) = \sum_n \frac{\langle \pi_a | \psi_n \rangle \langle \psi_n | \pi_b \rangle}{i\hbar\omega_{\nu} + \mu - \epsilon_n} = \int d\epsilon \frac{D_{a,b}(\epsilon)}{i\hbar\omega_{\nu} + \mu - \epsilon} \quad (\text{F.12})$$

Now we introduce a model density of states that is diagonal in the orbital index and which has a constant density of states of width  $W$  centered at the energy  $E$

$$D_{a,b}(\epsilon) = S_{a,b}^{-1} \frac{1}{W} \cdot \underbrace{\theta\left(\epsilon - \left(E - \frac{W}{2}\right)\right)}_{=1 \text{ above lower bound}} \cdot \underbrace{\theta\left(\left(E + \frac{W}{2}\right) - \epsilon\right)}_{=1 \text{ below upper bound}} \quad (\text{F.13})$$

The density of states is normalized so that the total weight of each index equals unity. The Heaviside function  $\theta(x)$  is defined such that it vanishes for negative arguments and is equal to one for positive arguments.

We calculate the green's function for  $\mu = 0$ . The relative position of chemical potential and the band can be changed easily by adapting the position of the band, i.e. by changing  $E$ . Keeping  $E$  flexible rather than  $\mu$  allow expand the argument to multi-band situations. The resulting Green's function has the form

$$\begin{aligned} G_{a,b}(i\omega_\nu) &= \int_{E-W/2}^{E+W/2} d\epsilon \frac{\frac{1}{W} S_{a,b}^{-1}}{i\hbar\omega_\nu - \epsilon} \\ &= \frac{S_{a,b}^{-1}}{W} \int_{E-W/2}^{E+W/2} d\epsilon \frac{-\epsilon - i\hbar\omega_\nu}{(\hbar\omega)^2 + \epsilon^2} \\ &= -\frac{S_{a,b}^{-1}}{W} \int_{E-W/2}^{E+W/2} d\epsilon \frac{\epsilon + i\hbar\omega_\nu}{(\hbar\omega)^2 + \epsilon^2} \\ &= -\frac{S_{a,b}^{-1}}{W} \int_{E-W/2}^{E+W/2} dx \frac{x + i\hbar\omega_\nu}{(\hbar\omega)^2 + x^2} \\ &= -\frac{S_{a,b}^{-1}}{W} \left\{ \int_{E-W/2}^{E+W/2} dx \frac{x}{(\hbar\omega)^2 + x^2} + i\hbar\omega_\nu \int_{E-W/2}^{E+W/2} dx \frac{1}{(\hbar\omega)^2 + x^2} \right\} \\ &= -\frac{S_{a,b}^{-1}}{W} \left\{ \frac{1}{2} \ln [(\hbar\omega)^2 + (E + W/2)^2] - \frac{1}{2} \ln [(\hbar\omega)^2 + (E - W/2)^2] \right. \\ &\quad \left. + i\hbar\omega_\nu \left[ \frac{1}{\hbar\omega} \operatorname{atan} \left( \frac{E + W/2}{\hbar\omega} \right) - \frac{1}{\hbar\omega} \operatorname{atan} \left( \frac{E - W/2}{\hbar\omega} \right) \right] \right\} \\ &= \frac{S_{a,b}^{-1}}{W} \left\{ +\frac{1}{2} \ln \left[ \frac{(\hbar\omega)^2 + (E - W/2)^2}{(\hbar\omega)^2 + (E + W/2)^2} \right] - i \left[ \operatorname{atan} \left( \frac{E + W/2}{\hbar\omega} \right) - \operatorname{atan} \left( \frac{E - W/2}{\hbar\omega} \right) \right] \right\} \end{aligned} \quad (\text{F.14})$$

We used the integral formula's from Bronstein<sup>1</sup>

The result as function of  $\omega$  is shown in Fig. F.1 The real part is symmetric about the origin, while the imaginary part is antisymmetric. This is consistent with  $\mathbf{G}(-i\omega_\nu) = \mathbf{G}^\dagger(-i\omega_\nu)$ .

The imaginary part exhibits a step at the origin from  $\pi/W$  to  $\pi/W$  for metallic systems, that is for a finite density of states at the origin. For an insulating system the imaginary part of the Green's function is smooth and has a finite slope at the origin.

<sup>1</sup>Integral formula's:

$$\begin{aligned} \int dx \frac{1}{a^2 + x^2} &= \frac{1}{a} \operatorname{atan} \left( \frac{x}{a} \right) \\ \int dx \frac{x}{a^2 + x^2} &= \frac{1}{2} \ln (a^2 + x^2) \end{aligned} \quad (\text{F.15})$$

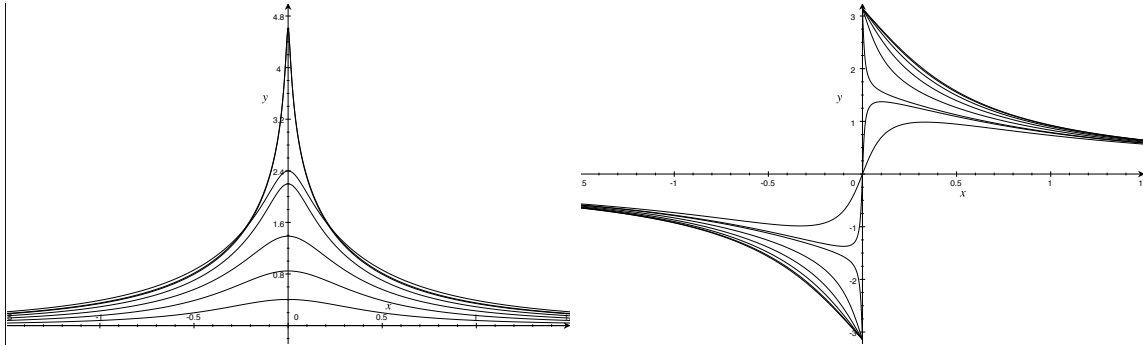


Fig. F.1: Negative real part (left) and negative imaginary part (right) of the Greens function with a constant density of states of width  $W = 1$ , centered at  $E = \{0, 0.1, 0.2, 0.3, 0.49, 0.51, 0.6\}$ .

The real part is positive if the band is centered above the chemical potential and it is negative if it is centered below. The real part vanishes if the density of states at the Fermi level. The real part becomes spiky exactly when the band touches the Fermi level. As the density of states shifts further to positive energies, the Greens function grows further in the tails, but it shrinks in the central region.

**The following is not correct! Please check!**

Limiting cases

- $\hbar\omega \rightarrow 0$ :



We use the expansion of the arcus tangens <sup>2</sup> and the expansion of the logarithm <sup>3</sup>.

$$\begin{aligned}
 G_{a,b}(i\omega_\nu) &= \frac{S_{a,b}^{-1}}{W} \left\{ +\frac{1}{2} \ln \left[ \frac{(\hbar\omega)^2 + (E - W/2)^2}{(\hbar\omega)^2 + (E + W/2)^2} \right] - i \left[ \text{atan} \left( \frac{E + W/2}{\hbar\omega} \right) - \text{atan} \left( \frac{E - W/2}{\hbar\omega} \right) \right] \right\} \\
 &\approx \frac{S_{a,b}^{-1}}{W} \left\{ +\frac{1}{2} \ln \left[ \left( \frac{E - W/2}{(E + W/2)} \right)^2 \right] + \frac{1}{2} \left[ \frac{1}{(E - W/2)^2} - \frac{1}{(E + W/2)^2} \right] \hbar\omega^2 \right. \\
 &\quad \left. - i \left[ \frac{\pi}{2} - \frac{\hbar\omega}{E + W/2} - \underbrace{\frac{\pi}{2} + \frac{\hbar\omega}{E - W/2}}_{\text{times -1 for insulators}} \right] \right\} \\
 &\approx \frac{S_{a,b}^{-1}}{W} \left\{ +\ln \left[ \frac{(E - W/2)}{(E + W/2)} \right] + \frac{EW}{\left( E^2 - \left( \frac{W}{2} \right)^2 \right)^2} \hbar\omega^2 - i \left\{ \begin{array}{ll} \frac{W}{E^2 - \left( \frac{W}{2} \right)^2} \hbar\omega & \text{for insulators} \\ \pm\pi - \frac{4|E|}{E^2 - \left( \frac{W}{2} \right)^2} \hbar\omega & \text{for metals} \end{array} \right\} \right\}
 \end{aligned} \tag{F.18}$$

•  $\hbar\omega \rightarrow \infty$ :

$$\begin{aligned}
 G_{a,b}(i\omega_\nu) &= \frac{S_{a,b}^{-1}}{W} \left\{ +\frac{1}{2} \ln \left[ \frac{(\hbar\omega)^2 + (E - W/2)^2}{(\hbar\omega)^2 + (E + W/2)^2} \right] - i \left[ \text{atan} \left( \frac{E + W/2}{\hbar\omega} \right) - \text{atan} \left( \frac{E - W/2}{\hbar\omega} \right) \right] \right\} \\
 &\approx \frac{S_{a,b}^{-1}}{W} \left\{ -EW(\hbar\omega)^{-2} - iW(\hbar\omega)^{-1} \right\} \\
 &\approx S_{a,b}^{-1} \left\{ -E(\hbar\omega)^{-2} - i(\hbar\omega)^{-1} \right\}
 \end{aligned} \tag{F.19}$$

---

2

$$\begin{aligned}
 \text{atan}(x) = y &\Rightarrow y = \tan(x) = \frac{\sin(x)}{\cos(x)} \\
 x = \frac{1}{\frac{\pi}{2} - y} + O(y - \frac{\pi}{2})^2 &\Rightarrow y \approx \frac{\pi}{2} - \frac{1}{x} \quad \text{for } x \rightarrow +\infty \\
 x = \frac{1}{y + \frac{\pi}{2}} + O(y + \frac{\pi}{2})^2 &\Rightarrow y \approx -\frac{\pi}{2} + \frac{1}{x} \quad \text{for } x \rightarrow -\infty
 \end{aligned} \tag{F.16}$$

3

$$\ln \left[ \frac{a+x}{b+x} \right] \approx \ln \left[ \frac{a}{b} \right] + \left( \frac{1}{a} + \frac{1}{b} \right) x + O(x^2) \tag{F.17}$$

# Appendix G

## Usage

### G.0.1 Control file

A typical control file looks as follows. The DMFT object is activated by the block !NTBO with the value MODUS='DMFT'. The DMFT interface requires a finite temperature calculation, which is specified by the !MERMIN block, where the temperature is specified. A finite temperature calculation requires SAFEORTHO=T, so that the wave function dynamics converges to eigenstates of the Hamiltonian.

```
!CONTROL
!GENERIC NSTEP=500 DT=5. START=F !END
!FOURIER EPWPSI=30. CDUAL=2.0 !END
!DFT TYPE=10
!NTBO MODUS='DMFT' !END
!END
!PSIDYN STOP=T FRIC=.05 SAFEORTHO=F
!AUTO FRIC(-)=0.3 FACT(-)=0.97 FRIC(+)=0.3 FACT(+)=1.0 MINFRIC=0.05 !END
!END
!MERMIN T[K]=4000. ADIABATIC=T RETARD=10. !END
!end
!EOB
```

### G.0.2 Structure file

A typical structure file may look as follows. New are the !NTBO subblocks.

```
!STRUCTURE
!GENERIC LUNIT[AA]= 3.8 EUNIT[EV]=T !END
!OCCUPATIONS EMPTY=10 NSPIN=2 SPIN[HBAR]=1.5 !END
!KPOINTS DIV=1 1 1 SHIFT=1 1 1 !END
!SPECIES NAME='CA' ID='CA_HBS_SC' NPRO=2 2 1
!NTBO NOFL=1 0 0 RAUG/RCOV=0.8 RTAIL/RCOV=1.4 TAILLAMBDA=2. 1.
CV=F FOCKSETUP=F !END
!END
!SPECIES NAME='MN' ID='MN_HBS' NPRO=1 1 1
```

---

```
!NTBO NOFL=1 0 1 RAUG/RCOV=1.2 RTAIL/RCOV=1.4 TAILLAMBDA=2. 1.
  LHFWeight=0.25 CV=F FOCKSETUP=F !END
!END
!SPECIES NAME='O_' ID='O_.75_6.0' NPRO=1 1 0
  !NTBO NOFL=1 1 RAUG/RCOV=1.2 RTAIL/RCOV=1.4 TAILLAMBDA=2. 1.
    CV=F FOCKSETUP=F !END
!END
!LATTICE T=  1.0000      0.0000      0.0000000000
            0.0000      1.0000      0.0000000000
            0.0000      0.0000      1.0000000000 !END
!ATOM NAME= 'CA1'  R=  0.0  0.0 0.0 !END
!ATOM NAME= 'MN1'  R=  0.5  0.5 0.5 !END
!ATOM NAME= 'O_1'  R=  0.0  0.5 0.5 !END
!ATOM NAME= 'O_2'  R=  0.5  0.0 0.5 !END
!ATOM NAME= 'O_3'  R=  0.5  0.5 0.0 !END
!ORBPOT_X
  !POT ATOM='MN(UP)1' VALUE=+.2 TYPE='D' RC=1.5 S=1 !END
!END
!END
!EOB
```

## **Appendix H**

# **Bloch representation**

## Appendix I

# Traditional formulation of the dynamical mean-field theory

The loop of solver takes as input a local Greens function and a U-tensor. It produces the value of the Luttinger-Ward functional, and its derivatives of the Luttinger-Ward functional with respect to Greens function and U-tensor.

The algorithm is the following

1. choose a hybridization function  $\Delta^{in}(i\omega_\nu)$
2. calculate an output Greens function  $\mathbf{G}^{out}$  as follows<sup>1</sup>: **Caution, I generalized Dieter's Equations without checking. Signs, factors, indices, etc. may be wrong!**

$$\begin{aligned}
 G_{\alpha,\beta}^{out}(i\omega_\nu) &:= -\frac{1}{Z} \int \prod_{\gamma} Dc_{\gamma}^* Dc_{\gamma} c_{\alpha}(i\omega_\nu) c_{\beta}^*(i\omega_\nu) e^{-S_{loc}} \\
 Z &= \int \prod_{\gamma} Dc_{\gamma}^* Dc_{\gamma} e^{-S_{loc}} \\
 S_{loc} &= -\int_0^{\beta} d\tau_1 \int_0^{\beta} d\tau_2 \sum_{\alpha,\beta} c_{\alpha}^*(\tau_1) \left[ i\hbar\omega_\nu \delta_{\alpha,\beta} - \Delta_{\alpha,\beta}^{in}(i\omega_\nu) \right] c_{\beta}(\tau_2) \\
 &\quad - \int_0^{\beta} d\tau \sum_{\alpha,\beta,\gamma,\delta} U_{\alpha,\beta,\delta,\gamma} c_{\alpha}^* c_{\beta}^* c_{\gamma} c_{\delta}
 \end{aligned}$$

Note, that only the hybridization function  $\Delta^{in}$  and the U-tensor enter in this calculation. The only result is the output Greens function  $\mathbf{G}^{out}$ .

3. Convert the Green's function into a self energy  $\Sigma^{out}$  using Dyson's equation

$$\Sigma_{\alpha,\beta}^{out}(i\omega_\nu) := -\left(G^{out}(i\omega_\nu)\right)_{\alpha,\beta}^{-1} + i\hbar\omega_\nu \delta_{\alpha,\beta} - \Delta_{\alpha,\beta}^{in}(i\omega_\nu) \quad (I.1)$$

---

<sup>1</sup>(D. Vollhardt *Dynamical Mean Field Theory for Strongly Correlated Materials* in *The LDA+DMFT Approach to Strongly Correlated Materials* E. Pavarini, E. Koch, D. Vollhardt and A. Lichtenstein Eds, Forschungszentrum Juelich 2011. Eq. 23.ff)

4. Use the Greens function  $\mathbf{G}$  passed to the solver as input argument (not  $\mathbf{G}^{out}$ !) to extract a new hybridization function  $\Delta^{out}$

$$\Delta_{\alpha,\beta}^{out}(i\omega_\nu) := -\left(G(i\omega_\nu)\right)_{\alpha,\beta}^{-1} + i\hbar\omega_\nu\delta_{\alpha,\beta} - \Sigma_{\alpha,\beta}^{out}(i\omega_\nu) \quad (1.2)$$

The steps given above yield a unique mapping from  $\Delta^{in}$  to  $\Delta^{out}$ . In other words the output hybridization function is a unique functional of the input hybridization function, the input Greens function and the U-tensor.

$$\Delta^{out} = F[\Delta^{in}, \mathbf{G}, \mathbf{U}] \quad (1.3)$$

The mapping depends on exactly two quantities, namely the input Green's function and the U-tensor. Self-consistency yields the fixed point  $\Delta^{out} = \Delta^{in}$  of this mapping. At this fixed point, Eq. 1.2 and Eq. 1.1 yield

$$G_{\alpha,\beta}^{out}(i\omega_\nu) = G_{\alpha,\beta}(i\omega_\nu) \quad (1.4)$$

and

$$G_{\alpha,\beta}(i\omega_\nu) = \left[ i\hbar\omega_\nu \mathbf{1} - \Sigma^{out}(i\omega_\nu) - \Delta^{in}(i\omega_\nu) \right]_{\alpha,\beta}^{-1}. \quad (1.5)$$

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